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PAIRING IN HADRON STRUCTURE*

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ABSTRACT

A many-body approach to hadron structure is presented, in which we consider two parton species: spin-0 (b-partons), and spin-$\frac{1}{2}$ (f-partons). We extend a boson and a fermion pairing scheme for the b-, and f-partons respectively, into a Yang-Mills gauge theory. The main feature of this theory is that the gauge field is not identified with the usual gluon field variable in QCD. We study the confinement problem of the hadron constituents, and obtain, for low temperatures, partons that are confined by energy gaps. As the critical temperatures for the corresponding phase transitions are approached, the energy gap gradually disappears, and confinement is lost. The theory goes beyond the non-relativistic harmonic oscillator quark model, in the sense of giving physical reasons why a non-relativistic approximation is adequate in describing the internal dynamics of hadron structure.

1. INTRODUCTION

In answering the question, What is a hadron made of ?, a reasonable answer lies in the quark-parton model (QPM). Feynman's interpretation [1,2] of the earliest lepton-scattering experiments [3] was in terms of elementary constituents, which he called partons. Later on these particles were identified with quarks and gluons, giving rise to the QPM. The gauge theory which takes these ideas into account, and has wide acceptation is Quantum Chromodynamics QCD [4].

However, the powerful gauge principle provides a convenient starting point for explaining the existence of partons; their forces may be derived from a unified point of view. This has led to the recent popularity of Grand Unified Theories, which attempt to encompass QCD, as well as other fundamental interactions. These efforts have somewhat shifted the elementary particle description from the quark level to a more elementary "preon level". So it might be useful to enquire, at this deeper level, whether a somewhat simpler picture of hadron structure might be possible. We will attempt to answer the above query in three stages:

(a) Pairing and the energy gap.

We might begin by looking at nuclear matter in order to obtain a first hint; nuclear structure exhibits many similarities with the electron structure of metals and alloys [5]. Bardeen, Cooper and Schrieffer (BCS), describe superconductivity in terms of pairing arising from residual interactions between the conduction electrons of the solid [6]. This, in turn, yields a good picture of the low energy exitations of the solid, which has the outstanding characteristic of having a gap in the energy spectrum. The BCS theory requires few changes to be in a position to describe many aspects of the nucleus [7], where the energy gap also appears. Beyond nuclear **physics**, in the realm of elementary particle physics, the analogy with the pairing phenomenon in solids has also been fruitful [8-12] ; in some of these applications the energy gap is of vital importance.

Another feature of pairing is its <u>universality</u>: besides superconductivity in metals, and pairing in the nucleus, this basic phenomenon also appears in the three superfluid phases of helium-3 [13]; in the interior of pulsars [14] ; possibly even in unusual systems, such as spin-aligned hydrogen, and helium-6 [15], and in the cosmic neutrino background [16]; although, in the latter case, the weak interactions will probably not overcome other mechanical effects opposing it. Therefore, in view of this universality, we might explore the possibility of its occurence amongst the preons , which through the pairing binding forces F_b , form the hadronic composite system, or mixture.

To universality, we would like to add a recent remark due to Feynman [17] : In a pure gauge theory of the type of Yang and Mills (YM) [18] , excluding quarks and leaving the gluons alone, it may be shown that an energy gap exists in the gluonic excitation spectrum (this general feature will probably not be affected by details in ref. [17] , which require some analytic reconsideration [19]). In other words, because of space correlation,

there are no long-wave excitations of arbitrarily low energy. Confinement is inferred appealing to the expectation for a Wilson loop [20], which falls exponentially with loop area, implying confinement. From these arguments we extract a hint as to how one may formulate a simpler theory of hadron structure:

(I) The lack of long-wave excitations of arbitrarily low energy may be related to confinement.

Yet, our approach towards the way in which our partons are confined, will not dwell on the detailed analysis of studying the space correlations which inhibit the low energy excitations. In fact, we shall guess the correct form of the ground state, the gap arising, as in other pairing theories, due to a highly correlated ground state. The best one can do here is to guess the correct form for the ground state, in the context of pairing; since no perturbation theory can lead to the gap, as it is functionally related to the coupling constant through an essential singularity (cf. Eqn. (2.33), below).

(b) Pairing and scalar partons.

We have only referred to pairing, thus far, amongst fermions. Scalar partons, on the other hand, may not be ruled out on purely experimental or theoretical grounds. Let us consider these two aspects carefully:

(i) By looking at the most recent data obtained from deep inelastic scattering of leptons off hadrons, we find that at finite energies, one could interpret the data in terms of hadron constituents which may be taken to be spin-$\frac{1}{2}$ partons (fermionic partons, or simply f-partons), and spin-0 partons (b-partons). This follows trivially from a consideration of the structure functions $F_1(x,Q^2)$, and $F_2(x,Q^2)$ of e-p scattering, for example. Here, x denotes the Feynman variable [1], and Q^2 is the invariant four-momentum transfer squared. It may be shown that for elastic scattering from point-like f-partons [21],

$$2x\, F_1(x) = F_2(x) \tag{1.1}$$

It is also useful to introduce the R-parameter, defined by,

$$R = \frac{F_2}{2x\, F_1}\left(1 + \frac{Q^2}{\nu^2}\right) - 1 \tag{1.2}$$

where ν is the energy lost by the lepton, and transferred to the hadrons. At finite energies $R \neq 0$. Yet, R is sensitive to the spin of the constituents [22] : For hadrons composed of f-partons alone, R = 0, whereas for partons of the b-type only, $R = \infty$. Hence, one may conclude that for a mixture of f- and b-partons, R falls between these values. In fact [22], $R = 0.21 \pm 0.1$, for leptons off protons, whereas $R = 0.24 \pm 0.1$, for leptons off deuterium.

(ii) In the unified gauge theory of Pati and Salam [23,24], scalar fields (b-partons) have been shown to play a leading role in spontaneously breaking the colour gauge symmetry, thereby providing mass to the gluons. Craigie and Salam [25], later showed that these scalar partons may be used to estimate deparures from standard QCD, particularly the R-parameter (1.2), which they find to be 0.28 as $x \to 0$; this should be compared with the experimental value of 0.24 from the data of leptons off deuterium, for example, as well as with the QCD prediction of 0.1 at small x.

Remarks (i) and (ii) lead us to identify a second important hint for a simpler formulation of hadron structure:

(II) The nucleon is a mixture of approximately 80% f-partons, and 20% b-partons.

In ref. [26], a pairing model of f-partons was studied, in which pairs were tentatively identified with b-partons, and the structure functions F_1 and F_2 were studied numerically. As opposed to such an example, we now put our main emphasis on pairing, not only for f-partons, but for b-partons as well, for all the reasons mentioned previously. This hope may be fulfilled, since it is known that pairing theories may be developed for bosons, as was originally shown by Valatin and Butler (VB) [27], and studied by others [28-31] ; all of these authors took advantage of the analogy with superconductivity, observed in liquid helium II. Yet, the energy gap, which turned out to be a difficulty in the VB model, will be an asset in the context of the hadronic mixture that concerns us here, since we wish to have an intrinsic effective mechanism for confinement. It should be noticed that these pairing theories do not require the "glue" as a separate field variable, but rather the formalism requires an averaged interaction of the usual Hartree-Fock type.

(c) Pairing and non-relativistic aspects of hadron structure.

A remarkable aspect of the spectroscopic predictions of the non-relativistic harmonic oscillator quark model is that it works so well in organizing the data, in the resonance region, for high energy phenomena; an intrinsic explanation for this feature is difficult [32].

On the other hand, our approach to the internal dynamics of the paired constituents of the hadron mixture, provides a clear physical argument indicating that the internal motion, in the rest frame of the mixture itself, is non-relativistic. Although we shall touch on this point in Section 4, we provide a brief account here: The pairing of f-partons is due to a binding force F_b , whose range $r > (1 \text{ Tev})^{-1}$, since the partons are preons, rather than quarks. At the same time the gap is not too large, since the constituents are mostly spin-$\frac{1}{2}$, as seen by deep inelastic probes of Q^2 of the order of a few (Gev)^2. Indeed, the evidence mentioned above,(II),shows this clearly. However, if our theory shows us that the order of magnitude of the fastest moving particles is given by

$$(\text{III}) \qquad \Delta\, r\,(F_b) \sim v_F$$

as will be shown in Section 4, the non-relativistic aspect of the f-parton motion is revealed. Here, v_F is the Fermi velocity, and Δ denotes the gap. The order of magnitude estimates for Δ and r indicate that $v_F << c$, the velocity of light. Thus, we have learned that the bound f-parton pair, although it has a wave function of very narrow spread, requires very little energy to excite a broken pair and, therefore,does not simulate an effective boson, leading us to the conclusion that both the gap and the Fermi velocity are small. The reason for the b-sector to be essentially non-relativistic is somewhat different: The VB interaction acts only between pairs of b-partons of vanishing total momentum p. We do not require each individual b-parton to be non-relativistic, but the bound

pair may be approximately described non-relativistically, since $p \gtrsim 0$ (in ref. [16] massive paired neutrinos were also described by a non-relativistic formalism, as the one in this paper).

The rest of the work will be distributed as follows: In Section 2 we give a brief account of the uncoupled pairing theories for the b- and f-partons, respectively. In Section 3, we transform the global invariance of the uncoupled theories into a local invariance, thereby coupling together the hadronic mixture, by means of a YM field. In Section4, we consider some qualitative features of the YM theory, including properties (I -III). Finally, in Section 5 we make a general summary, and arrive at some conclusions.

2. UNCOUPLED BINARY MIXTURE OF b- AND f-PARTONS

As in the usual theory of Quantum Liquids and Solids, we suppose that there exists a two-body interaction potential between the b-partons, which we denote by V (x - x′), and a two-body interaction potential between the f-partons, which we denote by U (x - x′).

We work in terms of a temperature -dependent hamiltonian H_B adjusted to the chemical potential λ_B since we shall consider states with an undetermined number of b particles. Hence, the role of the hamiltonian will be played by the operator ($H_B - {}_B N_B$), where N_B is the particle number operator [27, 31],

$$H_B - \lambda_B N_B = T_B + V \tag{2.1}$$

$$T_B = \sum_k (\varepsilon_k - \lambda_B) a_k^\dagger a_k \, , \quad \varepsilon_k = \frac{1}{2m_B} k^2 \tag{2.2}$$

$$V = \frac{1}{2} \sum_{kpq} V_k \, a^\dagger_{k+q} \, a^\dagger_{p-k} \, a_p \, a_q \tag{2.3}$$

where we have taken units such that $\hbar = 1$; m_B is the b-parton mass. Here, we have avoided the VB convention, and assume instead that $a_k^\dagger$ is a b-parton creation operator, V_k is the Fourier transform of the interparticle potential.

The hamiltonian of the f-parton system is similarly written down as,

$$H_F - \lambda_F N_F = T_F + U \tag{2.4}$$

$$T_F = \sum_k (\epsilon_k - \lambda_F) b_k^\dagger b_k \, , \quad \epsilon_k = \frac{1}{2m_F} k^2 \tag{2.5}$$

$$U = \frac{1}{2} \sum_{kpq} U_k \, b^\dagger_{q+k} \, b^\dagger_{p-k} \, b_p \, b_q \tag{2.6}$$

where the f-parton mass is denoted by m_F , and $b_k^\dagger$ is the f-parton creation operator.

We next find quasiparticle operators α_k, $\alpha_k^\dagger$, β_k, and $\beta_k^\dagger$ given by the Valatin-Butler [27], and Valatin Bogoliubov transformations [33, 34]:

$$\alpha_k = u_k^B a_k + v_k^B a^\dagger_{-k} \tag{2.7}$$

$$\beta_k = u_k^F b_k - v_k^F b^\dagger_{-k} \tag{2.8}$$

where, for canonical transformations,

$$(u_k^B)^2 - (v_k^B)^2 = 1 \tag{2.9}$$

$$(u_k^F)^2 + (v_k^F)^2 = 1 \tag{2.10}$$

the quantities $u_k^{B,F}$ and $v_k^{B,F}$ are c-numbers assumed real and even in k. For simplicity we only derive the gap equation for b-partons, and infer by analogy the corresponding one for f-partons. Whenever,

$$[H_B , \alpha_k^\dagger] = E_k^B \alpha_k^\dagger \tag{2.11}$$

is satisfied, we think that α_k generates an excitation of energy E_k^B from the ground state. The Eqn. of motion for α_k is,

$$[H_B , a_k] = -(\varepsilon_k - \lambda_B) a_k - \sum_{pq} V_{k-q} \, a^\dagger_{p-k} \, a_q a_p \tag{2.12}$$

In order to linearize these equations, we appeal to the VB transformation, and demand that the α's generate approximately independent excitations for different values of k. In this manner we are led to the thermal averages,

$$\langle a_k^\dagger a_{k'} \rangle = \delta_{kk'} \left\{ \left[(u_k^B)^2 + (v_k^B)^2 \right] \langle \alpha_k^\dagger \alpha_k \rangle + (v_k^B)^2 \right\} \equiv n_k^B \tag{2.13}$$

$$\langle a_k a_{-k'} \rangle = -\delta_{kk'} u_k^B v_k^B [2\langle \alpha_k^\dagger \alpha_k \rangle + 1] \equiv \chi_k^B \tag{2.14}$$

Therefore, the linearized equation of motion becomes:

$$[H_B, a_k] = -(\varepsilon_k - \lambda_B) a_k - V_0 N_B a_k - \sum_q V_{k-q} (n_q^B a_q + \chi_q^B a_q^\dagger) \tag{2.15}$$

The most convenient form for diagonalizing Eqn. (2.15) is in terms of the following definitions

$$\Delta_k^B = \sum_{k'} V_{k-k'} \chi_{k'}^B \tag{2.16}$$

$$\xi_k^B = \sum_{k} V_{k-k'} n_{k'}^B \tag{2.17}$$

$$\tilde{\varepsilon}_k = \varepsilon_k - \lambda_B + V_0 N^B + \xi_k^B \tag{2.18}$$

In other words, the linearized equation of motion may be conveniently written as,

$$[H_B, a_k] = -\tilde{\varepsilon}_k a_k - \Delta_k^B a_{-k}^\dagger \tag{2.19}$$

Diagonalization is achieved with the VB transformation, so as to turn Eqn. (2.19) into the original Eqn. (2.11); this process allows us to write the following identities:

$$E_k^B = \sqrt{[\tilde{\varepsilon}_k^2 - (\Delta_k^B)^2]} \tag{2.20}$$

$$(u_k^B)^2 + (v_k^B)^2 = \tilde{\varepsilon}_k / E_k^B \tag{2.21}$$

$$u_k^B v_k^B = \Delta_k^B / 2E_k^B \tag{2.22}$$

However, for b-partons, we have,

$$\langle \alpha_k^\dagger \alpha_k \rangle = [\exp(\beta E_k^B) - 1], \ \beta = (k_B T)^{-1} \tag{2.23}$$

From the thermal averages $\langle a_k^\dagger a_{k'} \rangle$, $\langle a_k a_{-k'} \rangle$ and the identities (2.20) to (2.22), we infer,

$$n_k^B + 1/2 = \frac{\tilde{\varepsilon}_k}{2E_k^B} \coth(\beta E_k^B / 2) \tag{2.24}$$

$$\chi_k^B = -\frac{\Delta_k}{2E_k^B} \coth(\beta E_k^B / 2) \tag{2.25}$$

Finally, the gap equation follows from (2.25) and (2.16):

$$\Delta_k^B = -\sum_{k'} V_{kk'} \frac{\Delta_{k'}^B}{2E_{k'}^B} \coth(\beta E_{k'}^B / 2) \tag{2.26}$$

The uncoupled f-parton pairing theory is analogous; but we obtain a modified single-particle excitation energy from the ground state,

$$E_k^F = \sqrt{[\tilde{\epsilon}_k^2 + (\Delta_k^F)^2]} \tag{2.27}$$

as well as a modified gap equation,

$$\Delta_k^F = -\sum_{k'} U_{kk'} \frac{\Delta_{k'}^F}{2E_{k'}^F} \tanh\left(\beta E_{k'}^F/2\right) \qquad (2.28)$$

The gap equations may be solved for simple potentials, for example a popular choice is [6, 27]:

$$V_{kk'} = \begin{cases} V, & \text{for } |k|, |k'| \leq \delta \\ 0, & \text{for } |k|, |k'| > \delta \end{cases} \qquad (2.29)$$

$$U_{kk'} = \begin{cases} -U, & \text{for } |k|, |k'| \leq \delta \\ 0, & \text{for } |k|, |k'| > \delta \end{cases} \qquad (2.30)$$

where δ is some small parameter; in this case of purely repulsive interaction amongst the b-partons, and purely attractive amongst the f-partons, we may integrate the gap equations, and infer the corresponding critical temperature for a phase transition, from the conditions :

$$\Delta^F(T_c^F) = 0 \quad ; \quad \Delta^B(T_c^B) = 0 \qquad (2.31)$$

Then, comparing with the zero temperature limit, one is led to the relations between (0), and T_c , which we write symbollically as:

$$\Delta^{B,F}(0) \propto k_B T_c^{B,F} \qquad (2.32)$$

Here, k_B is the Boltzmann constant. The theory, in the uncoupled limit, is capable of predicting both the gap and the critical temperature. It is a strightforward exercise with conditions (2.30) to infer the well known BCS result:

$$\Delta^F(0) \sim 2\delta \exp\left(-1/\rho_F U\right) \qquad (2.33)$$

where the expression for fermion density ρ_F will be given in Section 3. A similar expression may be derived for $\Delta^B(0)$. However, since excited pairs of bound b-fermions are possible within the scope of this theory (cf. discussion in Section 4), a phenomenon analogous to "glue-ball" excitation occurs in the theory, in the sense that in such excitation no f-partons take part. The theory also predicts their mass as being of the order $\Delta^B(0)$. A good quantitative prediction will require a full solution of the gauge theory, which we proceed to develop in the next section.

3. INTERACTING BINARY MIXTURE OF b- AND f-PARTONS

The interaction between b- and f-partons will be achieved in the context of a gauge theory. For this purpose, a lagrangian formulation of the uncoupled pairing theories is required. The full c-number configuration space of superconductivity [35] is perhaps too complicated for our purpose, and the more accesible isospin formulation [36] is adequate. We may start from a model lagrangian density [37], which we shall modify for consistency with Section 2:

$$\begin{aligned} \mathcal{L}_F(\vec{x},t) = {} & i\,\psi_\sigma^\dagger(\vec{x},t)\,\partial_t\psi_\sigma(\vec{x},t) \\ & - \frac{1}{2m_F}\,\partial_{\vec{x}}\psi_\sigma^\dagger\cdot\partial_{\vec{x}}\psi_\sigma(\vec{x},t) - \lambda_F\,\psi_\sigma^\dagger(\vec{x},t)\psi_\sigma(\vec{x},t) \\ & + U\,\psi_\uparrow^\dagger(\vec{x},t)\psi_\downarrow^\dagger(\vec{x},t)\psi_\downarrow(\vec{x},t)\psi_\uparrow(\vec{x},t) \end{aligned} \qquad (3.1)$$

We assume, for the moment, that m_F is sufficiently large to allow this lagrangian density to be a good approximation, but we shall return to the validity of the non-relativistic approximation in the following section. Here we have preferred to measure the kinetic energy relative to the Fermi energy level. We have also preferred:

(i) to avoid the symbolic form for the kinetic energy term, namely $\psi_\sigma^\dagger \tilde{\epsilon}_p \psi_\sigma$ [37], since the expression adopted in (3.1) is clearer;

(ii) for simplicity we have adopted a purely attractive and local f-parton-f-parton interaction potential,

$$U(|\vec{x}-\vec{x}'|) = -U\delta^{(3)}(\vec{x}-\vec{x}') \qquad (3.2)$$

instead of the simple BCS type of selection (2.30);

(iii) to be consistent with our selection of hamiltonian (2.4) to (2.6), we have allowed for a chemical potential, unlike ref. [37].

We remark that such a lagrangian density is invariant under the global symmetry:

$$\psi_\sigma \to \psi_\sigma' = e^{i\alpha}\psi_\sigma \qquad (3.3)$$

where α is some constant. From Nöther's Theorem, we infer the corresponding current and charge density:

$$\vec{J}_F = \frac{-i}{2m_F}\sum_\sigma\left[\psi_\sigma^\dagger \partial_{\vec{x}}\psi_\sigma - (\partial_{\vec{x}}\psi_\sigma)^\dagger\psi_\sigma\right] \qquad (3.4)$$

$$\rho_F(\vec{x},t) = \sum_\sigma \psi_\sigma^\dagger(\vec{x},t)\psi_\sigma(\vec{x},t) \qquad (3.5)$$

These quantities satisfy the continuity equation:

$$\partial_{\vec{x}}\vec{J}_F + \partial_t \rho_F = 0 \qquad (3.6)$$

which, in turn, implies that the total number of f-partons

$$N_F = \int \rho_F(\vec{x},t)\, d\vec{x} \qquad (3.7)$$

is conserved in time.

We have to be more careful in the case of pairing of b-partons: in Section 2 we have retained the anomalous pair thermal averages $\langle a_k\, a_{-k}\rangle$ (cf. Eqn.2.14); in other words, the interaction that forms pairs is less general than the matter-field hamiltonian would allow; that is, the VB interaction includes only interaction between pairs of b-partons of zero total momentum. For this reason we express the total field operator ϕ(x, t) for the b-parton, in terms of the eigenfunctions of the free hamiltonian T_B (Eqn. 2.2), which form a complete orthonormal set, as follows:

$$\phi_+ = \sum_{k>0} a_k(t)\, u_k(\vec{x}) \qquad (3.8)$$

$$\phi_- = \sum_{k>0} a_{-k}(t)\, u_{-k}(\vec{x}) \qquad (3.9)$$

and consider ϕ_S, where S = +, or – . Hence, we may write the b-parton pairing theory in terms of the lagrangian density,

$$\mathcal{L}_B(\vec{x},t) = i\,\phi_S^{\dagger}(\vec{x},t)\,\partial_t \phi_S(\vec{x},t) - \frac{1}{2m_B}\partial_{\vec{x}}\phi_S^{\dagger}(\vec{x},t)\cdot\partial_{\vec{x}}\phi_S(\vec{x},t) - \lambda_B \phi_S^{\dagger}(\vec{x},t)\phi_S(\vec{x},t) - V\phi_+^{\dagger}(\vec{x},t)\phi_-^{\dagger}(\vec{x},t)\phi_-(\vec{x},t)\phi_+(\vec{x},t) \tag{3.10}$$

As in the case of $\mathcal{L}_F$, we assume that m_B is sufficiently large to allow this density to be a good approximation, but we shall return, in the next section to the validity of the non-relativistic approximation.

Once again we observe global invariance $\delta\phi_S = i\beta\phi_S$, from which we infer the corresponding current $\vec{J}_B$ and charge density $\rho_B(x, t)$, analogous to Eqns. (3.4) and (3.5); these quantities also satisfy the continuity equation,

$$\partial_{\vec{x}}\vec{J}_B + \partial_t \rho_B = 0 \tag{3.11}$$

and the total number of b-partons,

$$N_B = \int \rho_B(\vec{x},t)\,d\vec{x} \tag{3.12}$$

is conserved in time. It seems convenient to define the b-parton concentration parameter R^B :

$$R^B = N_B\left(N_B + N_F\right)^{-1} \tag{3.13}$$

since this is an experimentally accessible quantity; once we couple both species in the hadron mixture, Eqn. (3.13) will provide an interesting test for the theory (cf. condition II). Taking advantage of the"isotopic" symmetry:

$$\Psi(\vec{x},t) = \begin{pmatrix} \psi_{\uparrow}(\vec{x},t) \\ \psi_{\downarrow}(\vec{x},t) \end{pmatrix} \tag{3.14}$$

we identify an analogous symmetry for b-partons:

$$\Phi(\vec{x},t) = \begin{pmatrix} \phi_+(\vec{x},t) \\ \phi_-(\vec{x},t) \end{pmatrix} \tag{3.15}$$

and obtain, in an evident notation (τ_3 is Pauli's matrix):

$$\mathcal{L}_F = i\,\Psi^{\dagger}\partial_t\Psi - \frac{1}{2m_F}\partial_{\vec{x}}\Psi\cdot\tau_3\,\partial_{\vec{x}}\Psi - \lambda_F\Psi^{\dagger}\Psi + \frac{1}{2}U\,\Psi^{\dagger}\tau_3\Psi\Psi^{\dagger}\tau_3\Psi \tag{3.16}$$

In Eqn. (3.16) a c-number term , arising from a rearrangement of non-commuting operators, has been neglected, since this will not affect our results. Similarly, we obtain:

$$\mathcal{L}_B = i\,\Phi^{\dagger}\partial_t\Phi - \frac{1}{2m_B}\partial_{\vec{x}}\Phi^{\dagger}\cdot\tau_3\,\partial_{\vec{x}}\Phi - \lambda_B\Phi^{\dagger}\Phi - \frac{1}{2}V\Phi^{\dagger}\tau_3\Phi\Phi^{\dagger}\tau_3\Phi \tag{3.17}$$

We introduce interaction between the parton species in the hadronic mixture by turning the global symmetries:

$$\delta\Psi = i\varepsilon\,\tau_3\Psi \tag{3.18}$$

$$\delta\Phi = i\varepsilon'\,\tau_3\Phi \tag{3.19}$$

into the same local invariance, by introducing the same compensating Yang-Mills gauge field to enforce the local symmetry. The physical motivation for this assumption lies not only on simplicity, but it also relies on the (Abelian) gauge theory for He II [38-41]. In that theory the gauge field does not appear as the "glue" of the interaction. Through a Madelung transformation [42], such an Abelian gauge field appears as a velocity field, in a two-fluid hydrodynamics. In the present case, the underlying multifluid non-Abelian hydrodynamics will, under the corresponding Madelung transformation, yield a gauge field which will play the role of some collective effect; but for the first qualitative predictions of the theory, which we envisage in this work, we need not develop this point any further. Therefore, following the work of YM, we enforce the local symmetry by letting all derivatives of Ψ and Φ appear as the covariant derivatives,

$$\partial_t - ig\vec{B}_0 \cdot \vec{\tau} \tag{3.20}$$

$$\partial_{x_i} - ig\vec{B}_i \cdot \vec{\tau} \tag{3.21}$$

where $\vec{B}_\mu$ are (2x2) matrices which compensate for the loss of global symmetry as the constants α and β acquire coordinate dependence, turning into the functions $\mathcal{E}$ (x, t) and $\mathcal{E}'$ (x, t) of Eqns. (3.18) and (3.19). In this manner, we are led to the lagrangian density of the gauge theory for the interacting binary mixture of b- and f-partons:

$$\begin{aligned}\mathcal{L} = {} & i\Psi^{\dagger}(\partial_t - ig\vec{B}_0\cdot\vec{\tau})\Psi - \\ & -\frac{1}{2m_F}(\partial_{x_i} + ig\vec{B}_i\cdot\vec{\tau})\Psi^{\dagger}\cdot\tau_3(\partial_{x_i} - ig\vec{B}_i\cdot\vec{\tau})\Psi - \\ & -\lambda_F\Psi^{\dagger}\Psi + \frac{1}{2}U\Psi^{\dagger}\tau_3\Psi\Psi^{\dagger}\tau_3\Psi + \\ & + i\Phi^{\dagger}(\partial_t - ig\vec{B}_0\cdot\vec{\tau})\Phi - \\ & -\frac{1}{2m_B}(\partial_{x_i} + ig\vec{B}_i\cdot\vec{\tau})\Phi^{\dagger}\cdot\tau_3(\partial_{x_i} - ig\vec{B}_i\cdot\vec{\tau})\Phi - \\ & -\lambda_B\Phi^{\dagger}\Phi - \frac{1}{2}V\Phi^{\dagger}\tau_3\Phi\Phi^{\dagger}\tau_3\Phi + \frac{k}{4}\vec{F}_{\mu\nu}\cdot\vec{F}^{\mu\nu} - \\ & - k/2\,\vec{F}_{\mu\nu}\cdot(\partial^{\mu}\vec{B}^{\nu} - \partial^{\nu}\vec{B}^{\mu} + 2g\vec{B}^{\mu}\times\vec{B}^{\nu})\end{aligned} \tag{3.22}$$

where k is a dimensional parameter.

4. DISCUSSION

Some general features of the hadron mixture are apparent, once we have defined our lagrangian density without necessarily going into a full analytical solution of the Euler-Lagrange equations:

A clear qualitative prediction emerges, in the sense that for low T ($T \ll T_c^B$), the b-partons are confined, whereas for high $T \sim T_c^B$, the b-partons become free as broken pairs (BP); however, for intermediate temperatures excited pairs (EP) are a clear possibility [13]. Similarly, in the case of f-partons confinement occurs effectively at low temperatures; at intermediate temperatures, EP are also a possibility with masses $\gtrsim \Delta^F$ (the gap); confinement is lost as we approach T_c^F .

Since our partons are considered to be preons rather than quarks, the expected force binding the preons

$$F_b^{(Preon)} < 1/M_{quark}$$

On the other hand, the expected range of the binding force

$$r(F_b) > (1\ T_{ev})^{-1}$$

This gives us an indication of the order of magnitude of the coherence length ξ_o^F, which is linked to the gap by [6]:

$$\xi_o^F \cong V_F / \Delta^F(0) \qquad (4.1)$$

The magnitude of the coherence length is related to the size of the f-parton pair,

$$r(F_b) \sim \xi_o$$

Eqn. (4.1) may be written as our condition (III),

$$V_F \sim \xi_o^F \Delta^F(0) \qquad (4.2)$$

By the arguments of Section 1(c), we see that a small gap Δ^F(0) of a few Gev's, and a small coherence length, somewhat smaller than 10^{-3} fermis, imply

$$V_F < 10^{-3} \qquad (4.3)$$

therefore, providing an intrinsic explanation of why, for the internal dynamics of the hadron, in its rest frame, the non-relativistic generalized Hartree-Fock approximation presents us with a simple description of hadron structure. It is interesting to compare our work with earlier studies of non-relativistic internal hadron dynamics in its rest

frame, but the mixture itself undergoing relativistic motion [43].

5. CONCLUSIONS

The most exciting practical aspect of the theory is the relationship between the critical temperatures above which partons become unconfined, due to the onset of the corresponding phase transition. Relation between Δ and T_c is possible in the context of previous work [44, 17]. We feel that our work differs, in the sense that we have identified the hadron with a new form of the phenomenon of quantum liquids and solids. It lies beyond the scope of this work to explore all consequences of this identification, but we would like to conclude with the following remark: In other manifestations of quantum liquids, the highly correlated ground state may be destroyed not only thermally, but mechanically as well. For example, the gap disappears in superconductivity by the application of a sufficiently strong magnetic field $H > H_{c2}$, whereas superfluidity in He II disappears by the action of sufficiently high angular momentum. This presents us with alternative mechanisms for quenching the gap, and thereby, for the disappearance of the hadronic mixture as such. In particular, we would not expect to find hadrons of arbitrarily high (intrinsic) angular momentum.

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