In the above $p^0 = (\mathbf{p}^2 + m_K^2)^{1/2}$. Similarly for the state $|K^+(-\mathbf{p})\rangle$ let the quark momenta at rest be k_1' and k_2' which are Lorentz boosted by matrix $L(p')$ with $\mathbf{p}' = -\mathbf{p}$. We note that we have also $k_1'^0 = \lambda_1 m_K$ and $k_2'^0 = \lambda_2 m_K$. Let $G_E^{K1}(t)$ be the contribution for form factor where u-quark interacts. Since here $\bar{s}$ is spectator, we get momentum conservation equation as [19]

$$L(p)_{ij}k_2^j + L(p)_{i0}k_2^0 = L(p')_{ij}\mathbf{k}_2'^j + L(p')_{i0}k_2^0. \tag{5.13}$$

On multiplyng the above by the inverse of the 3×3 matrix $L(p)_{ij} = L(p')_{ij}$, we then obtain that

$$\mathbf{k}_2' = \mathbf{k}_2 + \lambda_2\frac{m_K}{p^0}\cdot 2\mathbf{p}, \qquad \text{or,} \qquad \mathbf{k}_1' = \mathbf{k}_1 - \lambda_2\frac{m_K}{p^0}\cdot 2\mathbf{p}. \tag{5.14}$$

We then replace $\mathbf{k}_1$ by the symmetric integration variable $\mathbf{k}$ with the substitutions

$$\mathbf{k}_1' = \mathbf{k} - \lambda_2\frac{m_K}{p^0}\mathbf{p}; \qquad \mathbf{k}_1 = \mathbf{k} + \lambda_2\frac{m_K}{p^0}\mathbf{p}. \tag{5.15}$$

Thus, when the u-quark interacts we obtain that

$$G_E^{K1}(t) = e_u\int \tilde{u}_K(\mathbf{k}_1')^\dagger\tilde{u}_K(\mathbf{k}_1)\left(f_1(\mathbf{k}_1')f_1(\mathbf{k}_1) + \mathbf{k}_1'\cdot\mathbf{k}_1 g_1(\mathbf{k}_1')g_1(\mathbf{k}_1)\right)d\mathbf{k}, \tag{5.16}$$

where the variables are as given in equation (5.15). We also note that spin rotations have been included [19], and, there is no contribution from the same as we have here $S(L(p'))^\dagger S(L(p)) = I$ for $\mathbf{p}' = -\mathbf{p}$.

There will be a parallel contribution $G_E^{K2}(t)$ where the $\bar{s}$ interacts and u is the spectator. This contribution is obtained in a similar manner as

$$G_E^{K2}(t) = -e_s\int \tilde{u}_K(\mathbf{k}_1')^\dagger\tilde{u}_K(\mathbf{k}_1)\left(f_3(\mathbf{k}_1')f_3(\mathbf{k}_1) + \mathbf{k}_1'\cdot\mathbf{k}_1 g_3(\mathbf{k}_1')g_3(\mathbf{k}_1)\right)d\mathbf{k}, \tag{5.17}$$

where $\lambda_2 \to \lambda_1$, and $e_u \to -e_s$. Here parallel to equation (5.15) we have $\mathbf{k}_1' = \mathbf{k} - \lambda_1\frac{m_K}{p^0}\cdot\mathbf{p}$, and, $\mathbf{k}_1 = \mathbf{k} + \lambda_1\frac{m_K}{p^0}\cdot\mathbf{p}$. On simplification we then obtain that

$$R_{chK}^2 = R_{chK1}^2 + R_{chK2}^2 \tag{5.18}$$

where, parallel to equation (5.7)

$$R^2_{chK1} = \left(\frac{2}{3}\lambda_2^2 + \frac{1}{3}\lambda_1^2\right) \times \int (\nabla \tilde{u}_K(\mathbf{k}))^2 d\mathbf{k} \tag{5.19}$$

is the contribution coming from the wave function of equation (3.15), and,

$$R^2_{chK2} = \frac{2}{3} \times \lambda_2^2 \int u_K(\mathbf{k})^2 \left\{ \frac{R_1^4 \mathbf{k}^2 \cos^2(\chi_1(\mathbf{k}))}{4(1-\cos^2\chi_1(\mathbf{k}))} + \frac{(1-\cos\chi_1(\mathbf{k}))}{\mathbf{k}^2} \right\} d\mathbf{k}$$
$$+\frac{1}{3} \times \lambda_1^2 \int u_K(\mathbf{k})^2 \left\{ \frac{R_3^4 \mathbf{k}^2 \cos^2(\chi_3(\mathbf{k}))}{4(1-\cos^2\chi_3(\mathbf{k}))} + \frac{(1-\cos\chi_3(\mathbf{k}))}{\mathbf{k}^2} \right\} d\mathbf{k}. \tag{5.20}$$

is the balance of the contribution. We may easily note that when $R_1 = R_3$ and $\lambda_1 = \lambda_2 = \frac{1}{2}$, the above expressions go over to the corresponding expressions for the pion. The first term in the curly brackets above came from the simplification

$$\cos\left(\frac{1}{2}\chi_1(\mathbf{k}_1')\right)\cos\left(\frac{1}{2}\chi_1(\mathbf{k}_1)\right) + \sin\left(\frac{1}{2}\chi_1(\mathbf{k}_1')\right)\sin\left(\frac{1}{2}\chi_1(\mathbf{k}_1)\right)$$
$$1 - \frac{R_1^4 \mathbf{k}^2 \cos^2\chi_1(\mathbf{k})}{6(1-\cos^2\chi_1(\mathbf{k}))} \times \lambda_2^2 \mathbf{p}^2, \tag{5.21}$$

and, thc second term came from

$$\sin^2\left(\frac{1}{2}\chi_1(\mathbf{k})\right)\hat{\mathbf{k}}_1' \cdot \hat{\mathbf{k}}_1 \to -\frac{2(1-\cos\chi_1(\mathbf{k}))}{3\mathbf{k}^2} \times \lambda_2^2 \mathbf{p}^2. \tag{5.22}$$

The above contributions are for $\bar{s}$ being the spectator, and, we have used equations (2.14) and (5.15) for the simplification. The second curly bracket arises on interchanging the two quarks.

We now have to estimate λ_1 and λ_2. We had earlier suggested [26] that for sharing of the energy at rest, the kinetic energies of the two constituents may be different, but the potential energy shall be equally shared. In the present determination of the mass of the kaon the potential picture is absent. We shall however extrapolate the same by looking at the expression in (4.6) to guess these factors. We shall consider here two possible identifications. From equation (4.6) let us "identify" the potential energy as

$$v_K = (2\pi)^3 \; N_K^2 \; \frac{1}{2}[m_1\langle vac'| - \bar{\psi}_3\psi_3|vac'\rangle + m_3\langle vac'| - \bar{\psi}_1\psi_1|vac'\rangle]. \tag{5.23}$$

The balance of the contributions contain only u terms or only s terms, which we identify as the respective kinetic contributions. We then obtain that [26] $\lambda_1 = 0.134$ and $\lambda_2 = 0.866$. This yields that $R^2_{chK1} = 4.39$ GeV^{-2} and $R^2_{chK2} = 1.20$ GeV^{-2}, so that

$$R_{chK} = 0.47 \text{ fms}. \tag{5.24}$$

We may otherwise identify that in equation (4.6) the $\langle uu \rangle$ part corresponds to $\lambda_1 m_K$, and, the $\langle \bar{s}s \rangle$ part corresponds to $\lambda_2 m_K$. We then have $\lambda_1 = 0.067$ and $\lambda_2 = 0.933$. This yields that $R^2_{chK1} = 5.04$ GeV^{-2} and $R^2_{chK2} = 1.37$ GeV^{-2}, so that

$$R_{chK} = 0.50 \text{ fms}. \tag{5.25}$$

The above values may be compared with the experimental value of $R_{chK} = 0.58$ fms [17]. The calculated value appears to be small, and indicates that taking the wave function as determined from *exact* chiral symmetry breaking may not be correct. The "identification" of the fractions λ_1, λ_2 is also unreliable. We *should* have a handle on spectroscopy which may clarify the above as well as give corrections to the kaon wave function as different from the purely vacuum structure contribution as in equation (3.15).

VI. Vacuum structure through a variational calculation

We shall now briefly note [10] how we can *derive* the above vacuum structure of chiral symmetry breaking from the QCD Lagrangian in the light quark sector. For this purpose as earlier [4,6,7] we shall use a variational ansatz for the ground state [10]. We shall then find that a constrained energy minimisation of the Hamiltonian leads to a QCD vacuum with both quark and gluon condensates for $\alpha_s > \alpha_c = 0.62$. Pion decay constant and the charge radius of the pion seem to fix the QCD coupling constant α_s as 1.28, which also yields that the bag pressure is given by $B_g^{1/4} \simeq 140$ MeV.

The QCD Lagrangian density is given as

$$\begin{aligned}\mathcal{L} = &-\frac{1}{2}G^{a\mu\nu}(\partial_\mu W^a{}_\nu - \partial_\nu W^a{}_\mu + g f^{abc} W^b{}_\mu W^c_\nu) + \frac{1}{4}G^a{}_{\mu\nu}G^{a\mu\nu}\\ &+\bar{\psi}(i\gamma^\mu\partial_\mu)\psi + g\bar{\psi}\gamma^\mu\frac{\lambda^a}{2}W^a_\mu\psi,\end{aligned} \tag{6.1}$$

where $W^a{}_\mu$ are the SU(3) colour gauge fields. We quantise in Coulomb gauge [27] and write the electric field $G^a{}_{0i}$ in terms of the transverse and longitudinal parts as

$$G^a{}_{0i} = {}^T G^a{}_{0i} + \partial_i f^a, \tag{6.2}$$

where f^a is to be determined. We take time t=0 [4] and use an expansion of the field operators $W_i^a(\mathbf{x})$, ${}^T G_{0i}^a(\mathbf{x})$ and $\psi(\mathbf{x})$ as earlier [4,6,7]. In Coulomb gauge, the expression for the Hamiltonian density, $\mathcal{T}^{00}$ from equation (6.1) is given as [28,29]

$$\begin{aligned}\mathcal{T}^{00} = &: \frac{1}{2}{}^T G^a{}_{0i}{}^T G^a{}_{0i} + \frac{1}{2} W^a{}_i(-\nabla^2) W^a{}_i + g f^{abc} W_i^a W_j^b \partial_i W_j^c \\ &+ \frac{g^2}{4} f^{abc} f^{aef} W_i^b W_j^c W_i^e W_j^f + \frac{1}{2}(\partial_i f^a)(\partial_i f^a) \\ &+ \bar\psi(-i\gamma^i\partial_i)\psi - g\bar\psi\gamma^i \frac{\lambda^a}{2} W_i^a \psi :,\end{aligned} \tag{6.3}$$

where : : denotes the normal ordering with respect to the perturbative vacuum, say $| \, vac >$. In order to solve for the operator f^a, we first note that

$$f^a = -W^a{}_0 - g\, f^{abc}\, (\nabla^2)^{-1} (W^b{}_i\, \partial_i W^c{}_0). \tag{6.4}$$

Proceeding as earlier [4] with a mean field type of approximation we obtain,

$$\begin{aligned}\nabla^2 W^a{}_0(\mathbf{x}) + g^2\, f^{abc} f^{cde} \, < vac' \,|\, W^b{}_i(\mathbf{x}) \partial_i (\nabla^2)^{-1} (W^d{}_j(\mathbf{x}) \,|\, vac' > \partial_j W^e{}_0(\mathbf{x})) \\ = J^a{}_0(\mathbf{x}),\end{aligned} \tag{6.5}$$

where,

$$J_0^a = g f^{abc} W_i^b {}^T G_{0i}^c - g\bar\psi\gamma^0 \frac{\lambda^a}{2}\psi. \tag{6.6}$$

As noted earlier [4,9], the solution of equation (6.5) does not suffer from Gribov ambiguity [30].

The solution for $W_0^a(\mathbf{x})$ will depend on the ansatz for the ground state $|vac' >$. The mean field type of solution in equation (6.5) does not have a perturbative analogy, and is not an exact solution of the problem. As earlier [4,9] we consider the same only for low energy phenomenology.

The trial variational ansatz for the QCD vacuum is now taken as [10]

$$|vac>= U_G U_F|vac> \tag{6.7}$$

obtained through the unitary operators U_G and U_F for gluons and quarks respectively on the perturbative vacuum $|vac>$.

For the gluon sector, the unitary operator U_G is of the form

$$U_G = \exp\left(B_G{}^{\dagger} - B_G\right) \tag{6.8}$$

with the gluon pair creation operator $B_G{}^{\dagger}$ given by [4]

$$B_G{}^{\dagger} = \frac{1}{2}\int f(\mathbf{k}) a^a{}_i(\mathbf{k})^{\dagger} a^a{}_i(-\mathbf{k})^{\dagger} d\mathbf{k} \tag{6.9}$$

In the above $a^a{}_i(\mathbf{k})^{\dagger}$ are the transverse gluon field creation operators satisfying the following quantum algebra in Coulomb gauge [10,29]

$$\left[a^a{}_i(\mathbf{k}), a^b{}_j(\mathbf{k}')^{\dagger}\right] = \delta^{ab}(\delta_{ij} - \frac{k_i k_j}{k^2})\delta(\mathbf{k} - \mathbf{k}') \tag{6.10}$$

with $a^a{}_i(\mathbf{k})$ annihilating the perturbative vacuum $|vac>$. Further $f(\mathbf{k})$ is a trial function associated with gluon condensates. Clearly Eqs. (6.8) and (6.9) correspond to operator equations which create an arbitrary number of gluon pairs. In fact Eq. (6.9) may be interpreted as an operator to create a Bose BCS state. We shall here take $f(\mathbf{k})$ to be spherically symmetric.

For the quark condensates, we use our earlier analysis of section II. Thus for the quark sector we take [6–8]

$$U_F = \exp\left(B_F{}^{\dagger} - B_F\right) \tag{6.11}$$

with

$$B_F{}^{\dagger} = \int \left[h(\mathbf{k}) q^i{}_I(\mathbf{k})^{\dagger}(\sigma\cdot\hat{k})\tilde{q}^i_I(-\mathbf{k})\right] d\mathbf{k} \tag{6.12}$$

This is the same as B_0 of equation (2.2) with $i = 1,2$. Now $h(\mathbf{k})$ is a *trial function* associated with quark antiquark condensates. We recall that the operators $q^{\dagger}$ and $\tilde{q}$ create quark and antiquark respectively when operating on the perturbative vacuum and satisfy usual anticommutation relations.

As already seen such a structure for the vacuum reduces to a Bogoliubov transformation for the operators. One can then calculate the energy-density functional given as

$$\epsilon_0 = < vac'|T^{00}(\mathbf{x})|vac' > \equiv \epsilon_0(h, f). \tag{6.13}$$

The expression for $\epsilon_0(h, f)$ as a functional of h and f is extremely complicated and has been explicitly noted in Ref. [10]. The condensate functions $f(\mathbf{k})$ and $h(\mathbf{k})$ are to be determined such that the above energy density is a minimum. Since these functions here cannot be determined analytically through functional minimisation, we choose the alternative approach [4,6] of parameterising the condensate functions.

For the gluon sector we take as earlier [4], with $k = |\mathbf{k}|$,

$$\sinh f(\mathbf{k}) = Ae^{-Bk^2/2} \tag{6.14}$$

This corresponds to taking a Gaussian distribution for the perturbative gluons in the nonperturbative vacuum [4].

For the quark condensate function we note that our phenomenological ansatz in section II corresponds to

$$\tan 2h(\mathbf{k}) = \frac{1}{(e^{R^2k^2} - 1)^{1/2}}. \tag{6.15}$$

We would like to generalise the above, so that like the gluon distribution function the above function can vanish more easily for all values of $\mathbf{k}$. We shall do so by now taking the ansatz as

$$\tan 2h(\mathbf{k}) = \frac{A'}{(e^{R^2k^2} - 1)^{1/2}} \tag{6.16}$$

where A' is a variational parameter. Considering the earlier success in the light quark sector, we would like to understand here through *theoretical considerations* how we can have $A' \simeq 1$. As earlier we shall relate the quark condensate function to the wave function of pion as a quark antiquark bound state [6–9]. We shall also choose the decay constant of pion as a constraint [10] during energy minimisation. We add here a remark regarding the flexibility of the ansatz of equation (6.16). We note that in this equation if we take $A' = Rm$ and keeping m fixed, let $R \to 0$, then this equation shall correspond to a free Dirac particle of mass m.

The energy density is then minimised [10] with respect to the condensate parameters subject to the constraints that the pion decay constant f_π and the gluon condensate value $\frac{\alpha_s}{\pi} < G^a_{\mu\nu}G^{a\mu\nu} >$ of Shifman Vainshtein and Zhakarov [3] come out as the experimental values of 92 MeV and 0.012 GeV^4 respectively. The result of such a minimisation shows the instability of the perturbative vacuum to formation of quark antiquark as well as gluon condensates when the coupling becomes greater than 0.62. Here $A' = A'_{min}$ is a function of $\alpha_s = g^2/4\pi$. We then note that for $\alpha_s = 1.28$, $A' \simeq 1$ and thus the ansatz of equation (6.16) becomes the same as that of section II. Thus for the above coupling constant the earlier calculations in the light quark sector as well as those for other hadronic properties [9] remain unchanged. We also then have $R \simeq 0.96$ fm. which is the correct length scale associated with QCD. The finite value of α_s corresponds to the low Q limit of $\alpha_s(Q)$ as obtained through optimised renormalisation group equations [31] where the calculated value was $\alpha_s \simeq 0.8$. We can also calculate here the bag constant as $B_g^{1/4} = (-\epsilon_0)^{1/4} \simeq 140$ MeV. We record that this constant has been evaluated through energy minimisation starting from the QCD Lagrangian of the light quark sector.

With the structure of QCD vacuum thus fixed from pionic properties and SVZ value we can also consider the meson correlators [11] as current-current correlators [12]. Consider the currents

$$J(x) = \bar{\psi}_i(x)\Gamma\psi_j(x); \quad \bar{J}(x) = \bar{\psi}_j(x)\Gamma'\psi_i(x). \tag{6.17}$$

In the above, Γ are matrices ($1, \gamma_5, \gamma_\mu$ or $\gamma_5\gamma_\mu$) and $\Gamma' = \gamma_0\Gamma^\dagger\gamma_0$, and, as earlier i, j are flavour indices. We then define the propagator for interacting quarks as [32]

$$S(x) = < vac'|T\psi(x)\bar{\psi}(0)|vac' > . \tag{6.18}$$

We can then show that [11]

$$R(x) = < vac'|T\ J(x)\bar{J}(0)|vac' > \equiv -Tr\left[S(x)\Gamma' S(-x)\Gamma\right]. \tag{6.19}$$

We also substitute $R_0(x) = -Tr\left[S_0(x)\Gamma' S_0(-x)\Gamma\right]$ for massless noninteracting quarks. After an evaluation of the above expressions [11] we estimate $R(x)/R_0(x)$ for currents corresponding to different meson channels, and compare the same with the calculations of others. The results are seen to be similar to those obtained earlier [12] through lattice gauge theory [33]. For pions however as usual it seems to be necessary to saturate the currents with intermediate pion states. The interesting fact here is that the present modelling contains a *microscopic* description of vacuum, which gets related to hadronic properties for phenomena [12] outside spectroscopy.

VII. Discussions

Let us recall what has been discussed here. We first neglect masses of the quarks and assume that global chiral symmetry breaks spontaneously.

This is described through a vacuum realignment where we approximate for the correlation of a quark at two different space points as in equations (2.13) and (2.14) with a Gaussian function. Thus for the vacuum structure for quark q_i, we introduce a single parameter R_i. This parameter also gives the four component quark field operators for exact chiral symmetry breaking. We then find that R_i can be determined from the experimental value of the decay constant. Further, for approximate chiral symmetry breaking, we obtain the masses of the mesons through current algebra as in equations like (4.6) or (4.8) in terms of the Lagrangian masses of the quarks. This introduces another constant m_i for each quark. With only these two parameters for each quark, we use the explicit form of vacuum realignment to draw conclusions for the hadronic properties and examine consistency of such a hypothesis.

For the light quark sector, we find that the parameter R for the vacuum structure as determined from f_π also yields R^2_{ch} correctly. For the charge radius, we use that mesons in motion should be obtained through Lorentz boosting [19]. It thus appears that *we know* the vacuum structure for quark condensates of the light quark sector from experimental observations of the above hadronic properties as conjectured earlier [9].

We also determine the vacuum structure of the s-quark sector from the experimental value of f_K, and, using the same, go on to derive the charge radius of the kaon. This falls short of the experimental value by about fifteen percent. We may recall that from chiral perturbation theory a similar disagreement is also there, where the theoretical value is larger by about the same amount [17].

We next calculate the decay constants of D and B mesons. In this sector a quantitative agreement of the same is not expected, but, we shall discuss about this again in the context of restoration of chiral symmetry. We find however that with the present hypothesis, the spinor structure for c and b quarks may be quite different from that of a free Dirac particle. This feature shall have consequences for spectroscopy of heavy mesons.

Let us now note some obvious limitations of the present calculations. The s-quark seems to have a Lagrangian mass of the order of 100 MeV, and the corresponding masses of c and b quark are higher. For them we have considered the wave functions of the mesons as obtained totally from chiral symmetry breaking. In some sense this may not be very bad as *post facto* the scales for chiral symmetry breaking for them are higher than the above masses. It is however desirable to look further into this. For this purpose it may be worthwhile to look at the heavy quark sector using ansatz of equation (6.16). This will increase the number of parameters, but also can thereby approximate real physics better. We may also include the current quark masses during extremisation, in which case the arithmetic shall be forbidding, but, we can proceed step by step. After all, it is in fact the current quark masses as parameters which *can* drive the condensate functions in different quark sectors to be different!

As $R_t \rightarrow \infty$, chiral symmetry gets restored. We note from equation (4.4) that the pion mass remains unaltered. From equation (3.12) however we note that then f_π continuously decreases as R^2 increases. We may recall that in equation (4.18) f_D as calculated is *larger* than the experimental upper bound of f_D of MARK-III [23]. This could also be an indication that when D-mesons are produced, there is some progress towards chiral symmetry restoration in u, d sectors, which decreases the value of f_D giving rise to a smaller upper bound against the zero temperature calculation. Also, f_{D_s} as observed in CLEO-III is smaller, but is closer to the experimental value, since chiral symmetry restoration in s-sector is likely to be slower. We note that J/ψ suppression [34] has been a conventional signature for QGP [35]. We could probably add to the same the dependance of the above decay constants as signal of *progress* towards chiral symmetry restoration in QGP.

The second part of the paper dealing with derivation of the above vacuum structure for the light quark sector only briefly recapitulates Ref. [10] along with some further results [11]. It is nice to see that we can successfully obtain the same, and that this seems to fix the value of $\alpha_s(Q)$ as $Q \rightarrow 0$ which corresponds to the finite value derived in Ref. [31].

The present work systemises Ref. [9] and then extends the same to s and other heavy quarks. It also gives a theoretical base for the derivation of such results [10], and, illustrates why the *ad hoc* assumptions of Ref. [19], applied to many coherent and incoherent processes [36,37], could be more successful than the earlier approaches to quark model.

It may be desirable to see how these results change with the present form of equation (2.12) related to vacuum structure of (2.14). As noted we may try a variational calculation with the ansatz of equation (2.14) replaced by equation (6.16), and we should also include the effects of the mass term of the Lagrangian to drive the condensates of the heavy quark sectors. The bright side here is that we now seem to know the vacuum structure for light quarks. However, it also emphasizes the limitations in our understanding the same for heavier quarks, while illustrating their relevance for the corresponding hadronic properties.

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