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Strange Quark, Heavy Quarks and Gluon Contents of Light Hadrons

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Abstract

Heavy-quark contents (such as $\bar{c}c$) and the gluon content in baryons as well as the $\bar{s}s$ content are calculated with the use of the trace anomaly of QCD and a chiral quark model where the constituent quark masses are generated by the spontaneous breaking of chiral symmetry. The result is applied to the examination of several quantities in detail such as the explicit symmetry breaking in baryons (e.g., the π -N sigma term), the intrinsic heavy quarks in baryons and the Higgs-baryon coupling constants. The empirical value of the sigma term is reproduced with a rather small $\bar{s}s$ content of the nucleon. The probability of finding charm in the nucleon is found to be 0.5%, which is consistent with the analysis of the charm structure function of the proton and is necessary to account for the J/ψ -production in proton-nucleous collisions. Both strange and heavy quarks turn out to be equally important for the Higgs-nucleon coupling. Small but non-negligible OZI-violation in baryons due to the axial anomaly plays a key role for all the above quantities. The heavy-quark and gluon contents of the pion and the kaon and the color-sextet quark content in nucleon are also discussed.

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1 Introduction

Effects of flavor mixing have been one of the main subjects of hadron physics. It is known that the mixing is small in the meson sector, which is summarized as the celebrated OZI-rule[1]; a few exceptions include the system of η and η' mesons where the axial anomaly gives a large violation of the OZI-rule. Recently there has arisen a growing interest in the OZI-violation processes in the baryon sector. The relevant observables are the πN sigma term ($\Sigma_{\pi N}$)[2], the quark-spin content of the proton i.e. the flavor singlet axial-charge g_A^0 of the nucleon[3], intrinsic charm quarks in the nucleon suggested by the J/ψ - and open-charm production in hadron-nucleus collisions [4], the strange-quark contribution to the Higgs-nucleon coupling [5, 6] and so on.

A purpose of this paper is to give a realistic estimate of the OZI-violation in baryons on the basis of an effective Lagrangian which can describe the OZI-violation in mesons quantitatively. A brief account of a part of this paper related to the light sea-quark content in baryons has been reported in [7]: this paper gives not only an extensive account of it but also presents an estimate of the heavy-quark and gluon contents of baryons, which are then applied to the examination of several observables intimately related with the violation of the OZI-rule in baryons.

Before going into detail, let us briefly summarize the reason why the observables we mentioned above are interesting. $\Sigma_{\pi N}$ is defined as the nucleon matrix element of the scalar operator

$$\Sigma_{\pi N} = \hat{m}(\bar{u}u + \bar{d}d)_N, \quad (1.1)$$

with $\hat{m} \equiv (m_u + m_d)/2$ being the average value of the current masses of u and d quarks. (We use the non-covariant normalization for the state vector $\langle p' | p \rangle = (2\pi)^3 \delta^3(\mathbf{p}' - \mathbf{p})$ throughout this paper unless otherwise stated.) Although the “empirical” value of $\Sigma_{\pi N}$ is still uncertain, the latest analysis shows $\Sigma_{\pi N} = 51 - 53$ MeV with an estimated error 6 – 8 MeV.[8] On the other hand, the 1st order chiral perturbation by the current quark masses

($m_{u,d,s}$) tells us that the ratio

$$y = \frac{2\langle ss \rangle_N}{\langle uu + dd \rangle_N}, \quad (1.2)$$

should be about 0.5 to reproduce the empirical value of $\Sigma_{\pi N}$ [2], which suggests a large OZI violation in the scalar channel in the nucleon.

The EMC data [3] for the spin dependent structure function of the proton suggest that the OZI-rule is largely violated in the axial vector channel, too;

$$\langle \bar{s}\gamma_\mu\gamma_5 s \rangle_N = (-0.19 \pm 0.06) \cdot \bar{N}\gamma_\mu\gamma_5 N, \quad (1.3)$$

at $Q^2 \simeq 10\text{GeV}^2$. ($\bar{N}\gamma_\mu\gamma_5 N$ is the normalization factor with the Dirac spinor N .) Although the problem is still controversial (see e.g. [9] and the reviews [10]), it might imply that there exists a sizeable s -quark content with the opposite helicity to the polarization of the proton.

The mass number (A) dependence $A^{c(x_F)}$ of the J/ψ -production in nuclei is known to have two components, which are dependent on the Feynman x_F . For small x_F , $\alpha(x_F) \sim 1$, i.e., the cross section behaves like A^1 which is consistent with the picture that production occurs through a hard process; another component behaves like $A^{2/3}$, which dominates A^1 -component for large $x_F \sim 1$. The latter component may be accounted for by the so-called diffractive dissociation in which the intrinsic *charm* content of the nucleon plays an essential role.[4] This intrinsic charm component which can be only described non-perturbatively should be attributed to an OZI-violating process in the nucleon wave-function and has several other implications.[11]

The Higgs-nucleon coupling constant is also sensitive to the sea-quarks in the nucleon. As was first pointed out by Shifman, Vainshtein and Zakharov (SVZ) [12], the coupling is mostly determined by the heavy sea-quarks (c, b, t). It was, however, subsequently questioned by Cheng[5] and Cheng[6] who pointed out that the strange-quark contribution due to the OZI-violation cannot be neglected if y is as large as 0.5 which enhances the coupling constant by a factor 2–3.

As can be seen in the above examples, it is important to estimate the heavy as well as the light quark content in baryons quantitatively. This is

possible for the quark content in the scalar channel, i.e., $\langle \bar{q}q \rangle_B >_B$ since one can use several sum rules such as the QCD trace anomaly [13] for the heavy quarks and the Feynman-Hellman theorem (FH) [14] for the light quarks, which directly relate the $\bar{q}q$ content to the above observables.

We will first discuss the light-quark content of baryons using the well-known FH theorem in sections 2 and 3;

$$\langle \bar{q}_i q_i \rangle_B = \frac{\partial M_B}{\partial m_i}, \quad (1.4)$$

where $i = (u, d, s)$, M_B is the baryon mass and m_i is the current quark mass of the i -th quark.³ To get M_B in terms of m_i , we take a chiral quark model that has (i) the constituent quark masses generated by the spontaneous breaking of chiral symmetry (SB χ S) and (ii) the confinement of quarks represented by the constituent quark model [15,16]. The generalized $SU_f(3)$ -Nambu-Jona-Lasinio (NJL) model [17] which incorporates the axial anomaly [18,19] will be adopted to describe SB χ S: the term responsible for the flavor mixing is the 't Hooft vertex induced by the axial anomaly [20,21,22].

Using the result of $\langle \bar{q}_i q_i \rangle_B$ calculated with (1.4), we will confirm that the small strangeness content ($y = 0.12$) and large $\Sigma_{\pi N} \simeq 50 \text{ MeV}$ can be obtained simultaneously [7], which was suggested by several authors [20,21,22,23]. The non-linear m_s dependence of the physical quantities [20,22,24] and the short-range interaction due to the one-gluon exchange [7] are essential to get the result.⁴ Then we will calculate the anomalous quark contents (i.e. the OZI violating contribution) in the other members of the octet and decouplet. In general, the anomalous quark contents due to the axial anomaly are small but considerable.

The heavy quark contents (c, b, t) of static baryons such as $\langle \bar{c}c \rangle_N$ are

³Note that if we use the covariant normalization for the state vector, the r.h.s. of (1.4) reads $\partial M_B / \partial m_i$.

⁴The effect of the one-gluon exchange on $\Sigma_{\pi N}$ was first considered by Nelson [25] using a naive constituent-quark model.

extracted from the trace anomaly in QCD;

$$M_B = \langle \Theta_\mu^\mu \rangle_B = \sum_q m_q (1 + \gamma_m) \langle \bar{q}q \rangle_B + \frac{\beta}{4\alpha_s} \langle G^2 \rangle_B, \quad (1.5)$$

where q runs for the heavy quarks ($q=h=c,b,t$) as well as light quarks ($q=l=u,d,s$). β is the Callan-Symanzik β -function and γ_m is the mass anomalous dimension. At the energy scale of order 1 GeV, γ_m is roughly 0.22 if one uses the one-loop result $\gamma_m = 2\alpha_s/\pi$ and $\alpha_s(1 \text{ GeV}) \simeq 0.35$. The effect of it can, however, be absorbed into a modification of the masses of the current quarks $m_q \rightarrow m_q^* \equiv m_q(1 + \gamma_m)$, so one may just regard it as $O(20\%)$ uncertainty of the current quark masses. We will evaluate the heavy-quark contents with the help of the heavy quark mass expansion [26] in section 4. It will be found that the intrinsic charm of the nucleon, for example, becomes

$$\frac{\langle \bar{c}c \rangle_N}{\langle \bar{u}u + \bar{d}d \rangle_N} \simeq 5 \times 10^{-3}, \quad (1.6)$$

which may be compared with the value (3×10^{-3}) expected from the analysis of the

charm structure function in the proton.[27] We also show that the heavy quark contents are larger in the decouplet baryons than those in the octet ones.

As a straightforward application of the calculation of the quark contents, we will evaluate the Higgs-baryon coupling constant $g_{\phi BB}$ in section 5. It will be shown that $g_{\phi BB}$ depends on the value of $\langle \bar{s}s \rangle_B$ and our result

$$g_{\phi NN} \simeq 1.2 \times 10^{-3} = 0.32 \cdot (\sqrt{2} G_F) \cdot M_N, \quad (1.7)$$

lies in the middle of the previous two extreme estimates using $y=0$ [12] and $y \simeq 0.5$ [5,6]. (Here G_F is the Fermi constant.)

Finally, we will briefly discuss the quark contents in the pion and the kaon using the NJL model in section 6. It will be proved that the heavy quark contents and the gluon content in the pion vanish in the chiral limit

$(m_{u,d} \rightarrow 0)$, which is most easily seen if one starts from the Gell-Mann-Oakes-Renner (GOR) relation [28]. The realistic values off the chiral limit will be found to be

$$\frac{\langle \bar{s}s \rangle_\pi}{\langle \bar{u}u + \bar{d}d \rangle_\pi} \simeq 1.3 \times 10^{-4} \quad \text{and} \quad \frac{\langle \bar{c}c \rangle_\pi}{\langle \bar{u}u + \bar{d}d \rangle_\pi} \simeq 3.0 \times 10^{-4}, \quad (1.8)$$

which are much smaller than the corresponding quantities for baryons. It is interesting that the pion-induced diffractive charm cross section also suggested an order of magnitude smaller value of the probability for intrinsic charm in pion compared with that in proton.[4] In our scheme, this difference is a consequence of the Nambu-Goldstone nature of pion. In the final section, we will give a brief summary and concluding remarks.

Here we should mention the scale dependence of the baryon matrix element of $\bar{q}q$. Since the operator is multiplicative renormalizable and the combination $m\bar{q}q$ is scale independent, we always take an optimal scale μ where $m(\mu)$ can be evaluated using the perturbation theory: the quantum fluctuations in the energy scale above μ are taken into account in $m(\mu)$ while those below μ are taken into account in the matrix element $\langle \bar{q}q \rangle_B(\mu)$. A similar division of the energy scale has been adopted in the evaluation of the hadron matrix element of the weak Hamiltonian with perturbative QCD corrections (see e.g. [29]). Following the strategy in that case, we will identify our cutoff of the effective theory with the optimum scale $\mu \simeq O(1)\text{GeV}$. The scale dependence of the quantity above $\mu = \Lambda$ is thus governed by the renormalization group, hence

$$\langle \bar{q}q \rangle_B |_\mu = \exp\left(\int_\Lambda^\mu \frac{\gamma_m}{\beta} dg\right) \cdot \langle \bar{q}q \rangle_B |_{\Lambda=1\text{GeV}}. \quad (1.9)$$

In the following, we do not explicitly write scales at which operators are normalized, but they should be always regarded at $\mu \simeq \Lambda \simeq 1\text{GeV}$.

2 The dynamical mass and the constituent quark model

As discussed in the introduction, the evaluation of the quark content of hadrons in the scalar channel reduces to obtaining the hadron masses as functions of the current quark masses.

The constituent quark model (CQM) [15,16] has great success in reproducing the mass spectrum of baryons. Therefore, it should be a good starting point for the study of the quark contents of baryons. An apparent difficulty in applying CQM to the problem of the quark-content is that the relation of the constituent quark masses $M_{u,d,s}$ with the current quark masses $m_{u,d,s}$ is unknown: $M_{u,d,s}$ are usually just fitting parameters in the model. This difficulty is overcome by taking the picture advocated by several authors in which $M_{u,d,s}$ are identified with the dynamical masses generated by the spontaneous breaking of chiral symmetry in QCD. (See e.g., [30,31,32,33].) Although one does not know the precise mechanism of the chiral symmetry breaking in QCD, it can be well described by the effective Lagrangian of the Nambu-Jona-Lasinio (NJL) type. [32,33]⁵

Following the idea of Goldman and Haymaker [31], let us decompose the low energy effective Lagrangian as follows;

$$\mathcal{L} = \mathcal{L}_{NJL} + \mathcal{L}_{OGE} + \mathcal{L}_{CONF}. \quad (2.1)$$

Here $\mathcal{L}_{NJL}$ is the NJL model to be specified later and is understood to be responsible for chiral symmetry breaking and flavor mixing. $\mathcal{L}_{OGE}$ and $\mathcal{L}_{CONF}$ denote the short-range one-gluon exchange interaction (OGE) and the force responsible for the confinement, respectively.

Our basic assumption is that the dynamics may be treated in two steps as suggested by Goldman and Haymaker. First $\mathcal{L}_{NJL}$ is turned on, hence the dynamical quark masses are generated. Then the short-range and long-range interactions represented by $\mathcal{L}_{OGE}$ and $\mathcal{L}_{CONF}$ act.

⁵A pedagogical review of the NJL model may be found in [34].

Now let us write down the explicit form of $\mathcal{L}_{NJL}$ [20,21];

$$\begin{aligned} \mathcal{L}_{NJL} = & \bar{q}(i\gamma \cdot \partial - \mathbf{m})q + \sum_{a=0}^8 \frac{g_s}{2} [(\bar{q}\lambda_a q)^2 + (\bar{q}i\lambda_a \gamma_5 q)^2] \quad (2.2) \\ & + g_D [\det \bar{q}_i (1 - \gamma_5) q_j + h.c.], \end{aligned}$$

where the quark field q_i has three colors ($N_c = 3$) and three flavors ($N_f = 3$), and λ_a ($a=0\sim 8$) are the Gell-Mann matrices with $\lambda_0 = \sqrt{\frac{2}{3}}\mathbf{1}$. The second term is the explicit $SU_V(3)(SU_f(3))$ -breaking part with $\mathbf{m} = \text{diag}(m_u, m_d, m_s)$ being the current quark mass matrix. As can be easily seen, the second term does not mix the flavors [35]; in other words, it represents an effective Lagrangian in the large N_c limit in QCD where all the virtual quark-loops can be neglected. Since the origin of such effective 4-ferm interaction has been discussed by many authors, we do not repeat it here and just quote some references [32,33]. The term responsible for the mixing of different flavors is thus the third term which represents the major non-leading term of the large N_c expansion i.e. the axial anomaly. This term has the $SU_c(3) \otimes SU_f(3)$ invariance but breaks the $U_A(1)$ symmetry. One can also regard this part as an effective interaction induced by the instanton [36]. The same kind of effective Lagrangian was first introduced by Kobayashi and Maskawa [18] as a clue to solve the η' -mass problem and has been examined in the long-wavelength limit by Schechter et al [19].⁶

Lagrangian (2.2) has several parameters, the current quark masses $\mathbf{m} = \text{diag}(m_u, m_d, m_s)$, the coupling constants g_s and g_D , and the momentum cutoff Λ characterising the scale of the chiral symmetry breaking. Our strategy is to fix these parameters in the meson sector and apply the same Lagrangian to the baryon sector.

For m_u and m_d , we assume the $SU_f(2)$ -invariance and take the standard value at 1GeV scale $\hat{m} = 5.5\text{MeV}$ [37]. Four other parameters were determined to reproduce the four basic quantities $m_\pi = 138\text{MeV}$, $f_\pi = 93\text{MeV}$, $m_K = 495\text{MeV}$ and $m_{\eta'} = 958\text{MeV}$. (g_s, Λ), m_s and g_D are fixed essentially by (m_π, f_π) , m_K and $m_{\eta'}$ respectively. The resulting parameter set becomes

⁶An $SU_f(2)$ -version of the 't Hooft vertex was extensively examined in [33] without recursing to the long-wave length limit.

[21]⁷

$$\Lambda = 631.4\text{MeV}, \quad g_s \Lambda^2 = 3.666, \quad g_D \Lambda^5 = -9.306, \quad m_s = 135.3\text{MeV}. \quad (2.3)$$

Here we have used a three-momentum cutoff scheme. If one uses the Lorentz-invariant cutoff, it has a value of $O(1\text{GeV})$ (see e.g. the detailed study [38]) which is identified with the typical scale of the chiral symmetry breaking and also the optimum scale μ mentioned before.

The basic quantity for our purpose is the dynamical quark masses which are obtained by solving the coupled gap equation[21];

$$\begin{aligned} M_u &= m_u - 2g_s \langle \bar{u}u \rangle_0 - 2g_D \langle \bar{d}d \rangle_0 \cdot \langle \bar{s}s \rangle_0, \\ M_d &= m_d - 2g_s \langle \bar{d}d \rangle_0 - 2g_D \langle \bar{s}s \rangle_0 \cdot \langle \bar{u}u \rangle_0, \\ M_s &= m_s - 2g_s \langle \bar{s}s \rangle_0 - 2g_D \langle \bar{u}u \rangle_0 \cdot \langle \bar{d}d \rangle_0, \end{aligned} \quad (2.4)$$

where $\langle \bar{q}q \rangle_0$ ($q=u,d,s$) is the i -th quark condensate

$$\begin{aligned} \langle \bar{q}q \rangle_0 &= -iN_c \text{tr} \int \frac{d^4 p}{(2\pi)^4} \frac{1}{\gamma \cdot p - M_i}, \\ &= -\frac{N_c}{\pi^2} \int_0^\Lambda p^2 dp \frac{M_i}{E_{ip}}, \\ &= -\frac{N_c \Lambda^3}{2\pi^2} \cdot x [\sqrt{1+x^2} - x^2 \ln(\frac{1+\sqrt{1+x^2}}{x})]_{x=M_i/\Lambda}, \end{aligned} \quad (2.5)$$

with $E_{ip} = \sqrt{M_i^2 + p^2}$. The solution of (2.4) using the parameter set (2.3) reads

$$M_u = M_d = 335\text{MeV} \quad \text{and} \quad M_s = 527\text{MeV}, \quad (2.6)$$

which are quite consistent with the phenomenological masses used in CQM [16]. The above coupled equation relates M_i with m_i in a non-linear way. As for $m_{u,d}$, the linear approximation of M_i is legitimate, but the legitimacy does not hold for m_s . In fact, it can be proved that there is a definite convergence radius ($\sim 50\text{MeV}$) for the expansion by m_i , which implies that one cannot expand the physical quantities by m_s ($= 130 - 170\text{ MeV}$) [39].

⁷We have made a finer tuning of the parameters than those made in [21], hence the parameters here are slightly different from those quoted there.

In table 1, we have listed the several physical quantities predicted by the model in the meson sector and their empirical values. The agreement is quite good in a 10 – 20% level. A few comments are in order here. (i) The calculation is in the tree level of the meson field, so we do not have chiral logs which in some cases have a sizeable effect. Such an example is the kaon decay constant which is known to be enhanced by 20% by the meson loop making our result in better agreement with the data.[41] (ii) In the meson sector, the confinement $\mathcal{L}_{CONF}$ has little effect, at least, on pion as is shown by Goldman and Haymaker [31]. The good agreement with the empirical values listed in Table 1 suggests the same situation holds even in the other mesons. (iii) The key parameter that controls the flavor mixing is g_D , i.e., the strength of the anomaly term in our Lagrangian. This is determined so as to reproduce the correct mixing angle in the $\eta - \eta'$ system and the mass of the η' -meson which are the most sensitive to the OZI-violation.

Now let us turn to the baryon sector by taking into account $\mathcal{L}_{OGE}$ and $\mathcal{L}_{CONF}$. The effects of confinement and OGE are summarized in the simple mass formula in CQM, e.g. by Isgur and Karl [16];

$$M_B = M_0 + \sum_i^{u,d,s} [M_i + \frac{a}{2M_i}] + b \sum_{i < j} \frac{\sigma_i \cdot \sigma_j}{M_i M_j}. \quad (2.7)$$

The meaning of the respective terms is as follows. M_0 represents the contributions of the confinement potential and the short-range color-electric interaction, which are independent of the constituent masses. $a/2M_i$ denotes the kinetic energy of the confined quarks. The spin-spin term with the coefficient b is the color-magnetic interaction responsible for the octet-decuplet mass splitting. Instead of taking the detailed form of the wave function of quarks inside the baryons to calculate the parameters M_0 , a and b , we just input the masses of proton(938), $\Delta(1232)$ and $\Omega(1672)$. Then we get

$$a = (172.3\text{MeV})^2, \quad M_0 = -52.9\text{MeV}, \quad b = (176.5\text{MeV})^3, \quad (2.8)$$

which gives the baryon masses as follows⁸:

$$\Lambda(1115) = 1114, \quad \Sigma(1193) = 1185, \quad \Sigma^*(1385) = 1372, \quad (2.9)$$

⁸The mass formula needed to obtain (2.8) and (2.9) are tabulated in ref.[7]

$$\Xi(1320) = 1331, \quad \Xi^*(1507) = 1519,$$

in MeV. The fits are excellent within the error of 1%. We note that a good fit is indispensable to extract the heavy quark contributions to be presented in the later section; if not one cannot get the reliable number for them.⁹

As we have fixed all the parameters of the model, we can calculate the quark contents without free parameters.

3 Light quark contents in baryons

Applying the FH theorem eq.(1.4) together with the mass formula (2.7), one gets the quark contents in baryons

$$\begin{aligned} \langle \bar{q}_i q_i \rangle_P &= 2(1 - \frac{a}{2M_u^2} + \frac{b}{M_u^3})Q_{ui} + (1 - \frac{a}{2M_u^2} + \frac{4b}{M_u^3})Q_{di}, \quad (3.1) \\ \langle \bar{q}_i q_i \rangle_{\Delta^{++}} &= 3(1 - \frac{a}{2M_u^2} - \frac{2b}{M_u^3})Q_{ui}, \\ \langle \bar{q}_i q_i \rangle_{\Delta^+} &= (1 - \frac{a}{2M_u^2} - \frac{2b}{M_u^3})(2Q_{ui} + Q_{di}), \\ \langle \bar{q}_i q_i \rangle_{\Sigma} &= [2 - \frac{a}{M_u^2} + 4b(\frac{M_u^2 M_s}{M_u^2 M_s} - \frac{1}{2M_u^3})]Q_{ui} + (1 - \frac{a}{2M_u^2} + \frac{4b}{M_u M_s^2})Q_{si}, \\ \langle \bar{q}_i q_i \rangle_{\Sigma^*} &= [2 - \frac{a}{M_u^2} - 2b(\frac{M_u^2 M_u}{M_u^2 M_u} + \frac{1}{M_u^3})]Q_{ui} + (1 - \frac{a}{2M_u^2} - \frac{2b}{M_u M_s^2})Q_{si}, \\ \langle \bar{q}_i q_i \rangle_{\Xi} &= [2 - \frac{a}{M_s^2} + 4b(\frac{M_u^2 M_u}{M_u^2 M_u} - \frac{1}{2M_s^3})]Q_{si} + (1 - \frac{a}{2M_u^2} + \frac{4b}{M_s M_u^2})Q_{ui}, \\ \langle \bar{q}_i q_i \rangle_{\Xi^*} &= [2 - \frac{a}{M_s^2} - 2b(\frac{M_u^2 M_u}{M_u^2 M_u} + \frac{1}{M_s^3})]Q_{si} + (1 - \frac{a}{2M_u^2} - \frac{2b}{M_s M_u^2})Q_{ui}, \\ \langle \bar{q}_i q_i \rangle_n &= 3(1 - \frac{a}{2M_u^2} - \frac{2b}{M_u^3})Q_{si}. \end{aligned}$$

⁹Here one should note that the instanton induced interaction in $\mathcal{L}_{NLL}$ can give the same kind of baryon mass splittings as those given by the conventional OGE [42]. Therefore we have not made a clear distinction between the two; rather b is interpreted to represent the sum of the effects coming from the two.

Here Q_{ji} is the i -th current-quark content in the j -th constituent-quark,

$$Q_{ji} \equiv \frac{\partial M_j}{\partial m_i}, \quad (3.2)$$

which is evaluated to be

$$Q_{ji} = \begin{pmatrix} 1 + 2g_s R_u & 2g_D (\bar{s}s)_0 R_d & 2g_D (\bar{d}d)_0 R_s \\ 2g_D (\bar{s}s)_0 R_u & 1 + 2g_s R_d & 2g_D (\bar{u}u)_0 R_s \\ 2g_D (\bar{d}d)_0 R_u & 2g_D (\bar{u}u)_0 R_d & 1 + 2g_s R_s \end{pmatrix}^{-1}, \quad (3.3)$$

with $R_i \equiv d < \bar{q}q_i >_0 / dM_i$. As is explicitly shown in Appendix A, Q_{ji} is found to be proportional to the scalar meson propagator at zero momentum transfer;

$$Q_{ji} \sim \int d^4x e^{-ikx} \langle T \bar{q}_j q_i(x) \bar{q}_i q_j(0) \rangle_0 |_{k^2 \rightarrow 0} \quad (3.4)$$

Accordingly Q_{ji} represents the effect of the scalar-meson exchange connecting the external scalar-operator $\bar{q}_i q_i$ with the j -th constituent quark in a baryon. Thus the OZI-violation in baryons in the scalar channel is partly related to the flavor-mixing in the propagator of the scalar mesons and is partly related to that in the coupling of the scalar-meson to the constituent quarks: Actually the scalar mesons are almost ideally mixed, thus one can only expect a small violation of the OZI-rule in the baryons in this channel. Furthermore, the off-diagonal terms of Q_{ji} are proportional to g_D , which clearly shows that the flavor mixing is due to the anomaly term in our model.

In Fig.1, a schematic figure of the light quark content in baryons is given. The wavy lines in the figure show the correlations between quarks by the introduction of $\mathcal{L}_{OGE,CONF}$, which are summarized in the M -dependent coefficients of Q_{ji} in eq.(3.1). The numerical value of $Q_{ji} \equiv (\mathbf{Q})_{ji}$ with the parameters in (2.3) reads

$$\mathbf{Q} = \begin{pmatrix} 2.14 & 0.45 & 0.15 \\ 0.45 & 2.14 & 0.15 \\ 0.28 & 0.28 & 1.42 \end{pmatrix}. \quad (3.5)$$

It is easily verified that the ratio y is not affected by the inclusion of $\mathcal{L}_{OGE}$ nor $\mathcal{L}_{CONF}$ and can be written only in terms of Q_{ji} :

$$y = \frac{2Q_{us}}{Q_{uu} + Q_{ud}}, \quad (3.6)$$

$$= \frac{-4g_D (\bar{d}d)_0 I_s^c(0)}{1 + 2g_s I_s^c(0)} \quad (3.7)$$

where $I_s^c(k^2)$ is the loop integral in the scalar channel due to the s -quark given in Appendix A. This equation clearly shows that y is proportional to the strength of the anomaly term.

One can see from eq.(3.1) that the gluon-exchange term has a positive contribution, though the kinetic term of the confined quark has a negative contribution. In general, if the baryon matrix element of an operator is negative (positive) and is proportional to a negative power of M_i , it contributes positively (negatively) to the quark contents. Because of this tendency, $\Sigma_{\pi N}$ is enhanced by $O(15\%)$ while the " $\pi - \Delta$ sigma term" $\Sigma_{\pi\Delta} \equiv \bar{m}(\bar{u}u + \bar{d}d)_\Delta$ is suppressed due to OGE.

The light quark contents in baryons are summarized in Table 2. Because of the large $m_s \sim 135\text{MeV}$, the OZI-violation in the nucleon induced by the axial anomaly (g_D) is rather small

$$y/2 \equiv \frac{(\bar{s}s)_N}{(\bar{u}u + \bar{d}d)_N} \simeq 0.06, \quad (3.8)$$

which should be compared with the value expected from the 1st order chiral perturbation ($y \simeq 0.5$). Note that the strength of g_D we have used is large enough to reproduce the η - η' mixing angle and the η' mass which crucially depend on the amount of the OZI-violation. On the other hand, $\Sigma_{\pi N}$ can have a large value even if y is small;

$$\Sigma_{\pi N} \simeq 50 \text{ MeV}. \quad (3.9)$$

A possible discrepancy with the empirical value ($\Sigma_{\pi N} = 51 - 53\text{MeV}$ with an estimated error 6-8 MeV [8]) could be attributed to the clouds of the pseudo-scalar(ps) mesons, which may give a further enhancement of $O(10)$

MeV [24,14]. Note that both the nonlinear effect of m_s (which was pointed out in ref.[24,22]) and the interaction between quarks (OGE) [7] are essentially important to give the large $\Sigma_{\pi N}$. Our conclusion here is also qualitatively consistent with the recent studies of the Skyrme model, which also yield a small y with the large $\Sigma_{\pi N}$. [23,43]

Although the $\bar{s}s$ content is smaller than $\bar{u}u$ or $\bar{d}d$ in nucleon, it is not negligible. Furthermore, the non-strange content in Ω^- is even larger than the strangeness content in the proton. This is simply because the $u(d)$ -quark has a mass (current and constituent) smaller than that of the s -quark, hence gets easily excited. For a detailed discussion of the anomalous quark content in other baryons, see the previous paper [7] by the present authors.

4 Heavy-quark and gluon contents in baryons

In the previous section, we have determined the light quark contents of baryons using CQM. Combining this dynamical calculation with the trace anomaly, one can extract the heavy-quark $h = (c, t, b)$ contents in baryons. For heavy quarks, the first-order quark mass expansion is supposed to be good [26];

$$\langle \bar{h}h \rangle_B \simeq -\frac{1}{12m_h} \langle \frac{\alpha_s}{\pi} G^2 \rangle_B, \quad (4.1)$$

which tells us that the heavy quark content scales as $1/m_h$ in the scalar channel. Inserting this into the trace anomaly equation (1.5), one gets the gluon content of baryons

$$\langle \frac{\alpha_s}{\pi} G^2 \rangle_B \simeq -\frac{9}{8} [M_B - \sum_i^{u,d,s} m_i \langle \bar{l}l \rangle_B], \quad (4.2)$$

and the heavy quark content in baryons

$$\langle \bar{h}h \rangle_B \simeq \frac{2}{27m_h} [M_B - \sum_i^{u,d,s} m_i \langle \bar{l}l \rangle_B]. \quad (4.3)$$

These equations were first derived in ref.[12] where light-quark contributions were assumed to be negligible on the r.h.s. of (4.2). In our case, we

have the privilege of already having calculated $\langle \bar{l}l \rangle_B$ in the previous section, thus we need not make such an assumption.

For the nucleon, one can further reduce the formula by using an approximate relation $\Sigma_{\pi N} \equiv \hat{m}(\bar{u}u + \bar{d}d)_N \simeq m_u(\bar{u}u)_N + m_d(\bar{d}d)_N$;

$$\langle \frac{\alpha_s}{\pi} G^2 \rangle_N \simeq -\frac{9}{8} [M_N - \Sigma_{\pi N} - m_s \langle \bar{s}s \rangle_N], \quad (4.4)$$

$$\langle \bar{h}h \rangle_N \simeq \frac{2}{27m_h} [M_N - \Sigma_{\pi N} - m_s \langle \bar{s}s \rangle_N], \quad (4.5)$$

where the empirically unknown parameter is only $\langle \bar{s}s \rangle_N$. These equations show that the theoretical values of the gluon and heavy-quark contents crucially depend on the strangeness content of the nucleon because m_s is large; for instance, if y were to be as large as 0.5, $\langle \bar{h}h \rangle_N$ would change almost by a factor 2 from the case of $y = 0$.

In Table 3, we present the gluon, the light-quarks and the charm contents of baryons. (We took $m_c(1\text{GeV})=1.35\text{GeV}$.) One can see that $\langle \bar{c}c \rangle_B$ does not vary so much from one baryon to another and takes ~ 0.05 . This could be expected since they are pure sea-quarks. On the other hand, the magnitude of the light-quark content depends on the baryons since they contain the valence contribution. If one looks at fine structures, one finds that the charm in the octet baryons is smaller than that of the decouplet;

$$\langle \bar{c}c \rangle_8 \simeq 0.05 \quad \text{and} \quad \langle \bar{c}c \rangle_{10} \simeq 0.07. \quad (4.6)$$

This is due to OGE; OGE enhances the light-quark contents in the octet and suppresses those in the decouplet as we saw in the previous section. Then the sea contributions have the opposite tendency since their sum is constrained by the sum rule. If one takes the ratio $\langle \bar{c}c \rangle_B / \langle \bar{l}l \rangle_B$, this tendency is further enhanced as is shown in the last column of Table 3, which means that the “probability” to find a charm quark is larger in the decouplet baryons.

Here some comments are in order on our result,

$$\frac{\langle \bar{c}c \rangle_N}{\langle \bar{l}l \rangle_N} \simeq 0.5\%. \quad (4.7)$$

Phenomenologically, the intrinsic charm in nucleon is suggested to be indispensable for the description of the charm-production in nuclei in terms of the diffractive dissociation [4]. Furthermore, the analysis of the charm structure function of the proton using the EMC data requires an intrinsic charm probability of order 0.3% [27]. However, non-perturbative estimates of the heavy quark content are scarce so far; the only paper that we are aware of is a calculation using the MIT bag model by Donoghue and Golowich [44]. They obtained the probability of finding charm content in the proton to be $\sim 2\%$. In their case, the charm admixture was solely attributed to the perturbative gluon exchange but with a large $\alpha_s \simeq 2.2$, so their result should be regarded as a rough estimate. In our case, on the other hand, the result is rather stable since the $\bar{c}c$ -content is determined by the hadron masses with a small admixture of the light quark contribution and seems consistent in magnitude with the phenomenological value cited above.

As for b and t quark contents, what one should do is to just multiply the mass suppression factor $m_c/m_{b,t}$ to eq.(4.7). Thus, for example,

$$\frac{\langle \bar{b}b \rangle_N}{\sum_i \langle \bar{l}l \rangle_N} \simeq 0.1\%, \quad (4.8)$$

for $m_b(1\text{GeV})=5\text{GeV}$. Recently, it is suggested that the bound states of the color-sextet quarks may be produced diffractively in high energy hadron collisions in the laboratory and may also be already observed in the cosmic ray experiments.[45] Since the dynamics of the sextet quark is governed by QCD [46], we can easily evaluate the sextet-quark (q_6) content of the baryons. If we take the optimum mass for $m_{q_6} \simeq 350\text{GeV}$ [47] determined to reproduce the mass of W and Z due to the $\langle \bar{q}_6 q_6 \rangle_0$ condensation, we get

$$\frac{\langle \bar{q}_6 q_6 \rangle_N}{\sum_i \langle \bar{l}l \rangle_N} \simeq 0.01\%, \quad (4.9)$$

where we have taken into account the enhancement factor 5 due to the higher representation. This number will help us to estimate the cross section to produce the sextet-quarks in the diffractive dissociation mechanism.

It is also interesting to note here that the value of $\langle \frac{2s}{\pi} G^2 \rangle_B$ is negative and $O(700 - 1000)\text{MeV}$ as is shown in the 2nd column of Table 3. QCD

sum rules tell us that the vacuum condensate $\langle \frac{2s}{\pi} G^2 \rangle_0$ is positive and is about $(330\text{MeV})^4$ [48]. By the comparison of the two quantities, the naive cleansing-out picture of the bag formation [49] seems to be valid for baryons;

$$\langle \frac{\alpha_s}{\pi} G^2 \rangle_B \simeq - \langle \frac{\alpha_s}{\pi} G^2 \rangle_0 \times \frac{4}{3} \pi R^3, \quad (4.10)$$

where $R \simeq 1\text{fm}$ is the typical size of the baryons (i.e. the bag radius) inside which the perturbative vacuum is realized.¹⁰ However one should take this result only as a qualitative one, since not only the vacuum energy difference but also the kinetic energy of valence quarks are included in the matrix element $\langle G^2 \rangle_B$.

5 Higgs-baryon coupling constants

Originally, the Higgs-nucleon coupling strength was derived using the trace anomaly combined with the heavy quark mass expansion by SVZ [12]: once one neglects the light quark contributions, the coupling constant is entirely determined by the heavy sea-quarks (c, b, t). This conclusion has, however, been questioned recently by Cheng[5] and Cheng [6], who pointed out that the strange-quark contribution cannot be neglected if one adopts $y \simeq 0.5$ and in fact it enhances the coupling strength by a factor 2–3. In this section, we will show that both pictures are not quite right but correspond to two extreme.

The simplest form of the coupling of the scalar Higgs-boson to quarks is written as

$$\mathcal{L}_{\phi qq} = \sum_q m_q (\phi/v) \bar{q}q, \quad (5.1)$$

with $v=(\sqrt{2}G_F)^{-1/2} = 246\text{GeV}$ being the symmetry breaking scale of the electroweak theory.

¹⁰The table also shows that the bag radii of the decoupled baryons are about 10% larger than those of the octet baryons.

The Higgs-baryon coupling constant is given by

$$\begin{aligned} v g_{\phi BB} &= \sum_i \langle m_i \bar{l} l \rangle_B + \sum_h \langle m_h \bar{h} h \rangle_B, \\ &\simeq \frac{2}{9} M_B + \frac{7}{9} \sum_i \langle m_i \bar{l} l \rangle_B, \end{aligned} \quad (5.2)$$

for 3 light and 3 heavy quarks.¹¹

From Table 3, we can read the light and heavy quark contributions to $v g_{\phi BB}$. For the nucleon, the light quark contribution reads

$$< m_u \bar{u} u + m_d \bar{d} d + m_s \bar{s} s \rangle_N \simeq 122 \text{ MeV}. \quad (5.3)$$

We remark that the magnitude of this quantity is correlated with that of y : if $y \simeq 0.47$ were adopted instead of our value $y \simeq 0.12$, the same quantity would become as large as 410 MeV. The heavy quark content reads

$$\sum_h m_h \langle \bar{h} h \rangle_N \simeq 181 \text{ MeV}. \quad (5.4)$$

Thus one sees that the light-quark contribution is not so large as Cheng [5] and Cheng [6] claim, but so small as SVZ [12] claims.

The coupling constants $v g_{\phi BB}$ as well as the light and heavy quark contributions to it and the ratio are tabulated in Table 4.

One finds that (i) the flavor mixing in the non-strange baryons generally increases $g_{\phi BB}$ since $\langle \bar{s} s \rangle_B$ is not negligible, and (ii) as the number of the valence s -quark increases, the Higgs-baryon coupling constant increases (the Higgs- Ω coupling is the largest) since $m_s \langle \bar{s} s \rangle_B$ dominates $m_u \langle \bar{u} u \rangle_B$ and $m_d \langle \bar{d} d \rangle_B$ owing to the large difference in the masses, $m_s/\bar{m} \simeq 25$.

Our value for the Higgs-nucleon coupling reads

$$v g_{\phi NN} = 303 \text{ MeV} = 0.32 M_N, \quad (5.5)$$

which is in the middle of the predictions $v g_{\phi NN} = 210 \text{ MeV} = 0.22 M_N$ ($y = 0$) [12] and $v g_{\phi NN} = 530 \text{ MeV} = 0.56 M_N$ ($y = 0.47$) [5, 6].

¹¹ Here the current quark masses are the ones which do not include the effect of the anomalous dimension.

An immediate application of this coupling constant is the constraint on the mass of the very light Higgs boson. (See the recent review, [50].) Although the combination of recent LEP data and low energy experiment excludes the possibility of the light Higgs boson in the standard model with the minimal Higgs sector, our calculation may be useful to give a constraint on the non-minimal Higgs models where the Higgs couplings to ZZ and $\bar{q}q$ are independent.

6 Quark- and gluon-contents of pion and kaon

In this section, we will examine the quark- and gluon-contents of the pion and the kaon using the FH theorem and the trace anomaly relation.

6.1 Quark and gluon contents of the pion

In applying the FH theorem, one needs to get the mass m_π in terms of the current quark masses m_i . This is nicely accomplished in the NJL model, where the meson masses are obtained from the dispersion relations the derivatives of which with respect to m_i can be performed straightforwardly. This direct procedure is certainly straightforward, but it is rather difficult to get the meaning of the terms obtained. On the other hand, as is well known, there is a approximate but nice formula for the pion mass i.e., the Gell-Mann-Oakes-Renner (GOR) relation [28] which enables us to get the intuition into the terms obtained:

$$m_\pi^2 = -\frac{\bar{m}}{f_\pi^2} (\bar{u}u + \bar{d}d)_0 + O(\bar{m}^2) + \dots \quad (6.1)$$

Here the pion decay constant f_π and the quark condensates are evaluated in the limit $m_{u,d} \rightarrow 0$ but not necessarily in the limit $m_s \rightarrow 0$; i.e., they can be a non-linear function of m_s . One can easily see this by looking at the derivation of the GOR relation carefully. (See e.g., ref [51].) We remark that the GOR formula is naturally derived from the dispersion equation in the NJL model [33] as it should be. Thus before the direct calculation, let us first start with the GOR formula.

The FH theorem applied to the pion reads

$$\frac{\partial m_\pi}{\partial m_s} = \langle \bar{q}q \rangle_\pi = \langle \bar{q}q \rangle_\pi^c / 2m_\pi, \quad (6.2)$$

where $\langle \cdot \rangle_\pi^c$ denotes the matrix element in the covariant normalization $\langle p' | p \rangle = 2\omega_p(2\pi)^3 \delta^3(\mathbf{p}' - \mathbf{p})$. Then inserting the GOR formula (6.1), one gets for the u and d contents

$$\langle \bar{u}u \rangle_\pi^c = \langle \bar{d}d \rangle_\pi^c = -\frac{1}{f_\pi^2} \langle \bar{u}u \rangle_0, \quad (6.3)$$

$$= \frac{m_\pi^2}{2\hat{m}}, \quad (6.4)$$

which could be also derived by using the soft pion theorem. As for the s content, we have

$$\langle \bar{s}s \rangle_\pi^c = m_\pi^2 \frac{\partial}{\partial m_s} \ln(|\langle \bar{u}u \rangle_0|/f_\pi^2), \quad (6.5)$$

which shows the remarkable fact that $\langle \bar{s}s \rangle_\pi^c$ vanishes in the chiral limit $m_{u,d} \rightarrow 0$. The r.h.s. of (6.5) gives $\langle \bar{s}s \rangle_\pi^c$ in the first order of the chiral perturbation, which can be calculated, for instance, with the use of the NJL model. (See Appendix B.)

Once we know the light quark content of the pion, the trace anomaly tells us the gluon and heavy quark contents: In fact, as for the baryons, we have

$$\langle \Theta_\mu^\mu \rangle_\pi^c = 2m_\pi^2, \quad (6.6)$$

$$\simeq m_u \langle \bar{u}u \rangle_\pi^c + m_d \langle \bar{d}d \rangle_\pi^c + m_s \langle \bar{s}s \rangle_\pi^c - \frac{9}{8} \langle \frac{\alpha_s}{\pi} G^2 \rangle_\pi^c, \quad (6.7)$$

where we have applied the heavy quark expansion in the second equality. Thus using eq.'s (6.1) and (6.3), we get the gluon content,

$$\langle \frac{\alpha_s}{\pi} G^2 \rangle_\pi^c \simeq \frac{8}{9} [m_\pi^2 + m_s \langle \bar{s}s \rangle_\pi^c], \quad (6.8)$$

and the heavy-quark content,

$$\langle \bar{h}h \rangle_\pi^c \simeq -\frac{2}{27m_h} [m_\pi^2 + m_s \langle \bar{s}s \rangle_\pi^c]. \quad (6.9)$$

Although the above derivation of the heavy-quark and gluon contents is similar with that adopted in the analysis of the Higgs decay into two pions [52], the present analysis of the $\bar{s}s$ content is quite new. It is worth mentioning here that the gluon and heavy-quark contents vanish in the chiral limit since $\langle \bar{s}s \rangle_\pi^c$ is proportional to the pion mass squared. This implies that only small values are allowed for the gluon and heavy-quark contents of the pion due to the chiral symmetry. On the other hand, such a constraint is absent in baryons. This difference between pion and nucleon seems to be consistent with the experimental data on the charm-production in hadron-nucleus collisions[4]; see below.

Now, let us calculate the quark contents of the pion by directly differentiating the pion mass obtained in the NJL model without recourse to the GOR relation which is valid only in the first order of the chiral perturbation. In the NJL model, the meson masses are obtained dynamically from the dispersion equations. The dispersion equation for the pion ($\pi^\pm$) reads;

$$0 = D_\pi^{-1}(m_\pi^2) \propto 1 + g_\pi \Pi_{\pi^\pm}(m_\pi^2), \quad (6.10)$$

where $g_\pi = g_\pi + g_D \langle \bar{s}s \rangle_0$ and $\Pi_{\pi^\pm}(k^2) = I_\pi^{\pi^\pm}(k^2) + I_\pi^{\pi^\pm}(k^2)$ with $I_\pi^{\pi^\pm}(k^2)$ being the loop integral in the pseudoscalar channel which is given in Appendix B. Differentiating the dispersion equation, one gets the r.h.s. of eq.(6.2) in terms of Q_{ji} since the loop integral is the function of the constituent quark mass M_i :

$$\langle \bar{q}q \rangle_{\pi^\pm}^c = \frac{1}{2J_u} [C_\pi(Q_{ui} + Q_{di}) + (\delta_{ui} + \delta_{di}) \frac{\langle \bar{u}u \rangle_0}{(M_u - \hat{m})^2}], \quad (6.11)$$

where

$$C_\pi = \frac{I_\pi^q(0)}{M_u - \hat{m}} - K_u - \frac{\langle \bar{u}u \rangle_0}{(M_u - \hat{m})^2}, \quad (6.12)$$

with

$$J_j = \frac{\partial I_j^P(m_\pi^2)}{\partial m_\pi^2}, \quad \text{and} \quad K_j = \frac{\partial I_j^P(m_\pi^2)}{\partial M_j}. \quad (6.13)$$

The explicit forms of J_j and K_j are given in Appendix B. We note that C_π which gives the strength of the flavor mixing in the pion can be easily

shown to be proportional to m_π^2 in the first order of the chiral perturbation and hence vanishes in the chiral limit. This is in accordance with the result given from the GOR relation.

Combining (6.11) with the trace anomaly, one can evaluate $\langle \bar{h}h \rangle_\pi$ and $\langle G^2 \rangle_\pi$. The numerical results without recourse to the chiral perturbation reads

$$\begin{aligned} \langle \bar{u}u \rangle_{\pi^\pm} &= \langle \bar{d}d \rangle_{\pi^\pm} = 6.2, \\ \langle \bar{s}s \rangle_{\pi^\pm} / \langle \bar{u}u + \bar{d}d \rangle_{\pi^\pm} &= 1.3 \times 10^{-4}, \\ \langle \bar{c}c \rangle_{\pi^\pm} / \langle \bar{u}u + \bar{d}d \rangle_{\pi^\pm} &= 3.0 \times 10^{-4}, \\ \langle \frac{\alpha_s}{\pi} G^2 \rangle_{\pi^\pm} &= -61.3 \text{ MeV}. \end{aligned} \quad (6.14)$$

Comparing (6.14) with (3.8) and (4.7), one sees that the sea-quarks in the pion are order of magnitude smaller than those in baryons, which is due to the Nambu-Goldstone nature of the pion. This may imply that the pion-induced diffractive charm cross section is order of magnitude smaller than the proton-induced cross section, which was in fact suggested by Brodsky, Peterson and Sakai [4] from the analysis of the experimental data. Another interesting point is that $\bar{c}c$ and $\bar{s}s$ contents in pion are comparable. This is not necessary against the intuition from the heavy quark mass expansion, since the expansion does not hold for the strange quark.

Using the quark contents obtained, we get the pion-Higgs coupling defined as the strength of the effective vertex $\mathcal{L}_{\text{int}} = g_{\pi\phi} \phi \vec{\pi} \cdot \vec{\pi}$;

$$g_{\pi\phi} = \sum_q m_q \langle \bar{q}q \rangle_\pi \cdot \frac{m_\pi}{v}, \quad (6.15)$$

$$= 84 \frac{m_\pi}{v} \text{ MeV}. \quad (6.16)$$

6.2 Quark and gluon contents of the kaon

As was done for the pion, we can evaluate the quark and gluon contents of the kaon from the knowledge of $\langle \bar{q}q \rangle_K = dm_K^2/dm_i$ ($i = u, d, s$) by utilizing the trace anomaly and the heavy-quark mass expansion.

One may, of course, start from the GOR relation for the kaon

$$m_K^2 = -\frac{1}{2}(m_u + m_s)\langle \bar{u}u + \bar{s}s \rangle_0, \quad (6.17)$$

which is given in the first order of the chiral perturbation for the s quark as well as the u quark. However, the deviation from the GOR relation for kaon is expected to be larger than that for pion, thus we will directly differentiate the dispersion equation of the kaon given in the NJL model:

$$1 + g_{K^\pm} \Pi_K(m_K^2) = 0. \quad (6.18)$$

Here $g_{K^\pm} = g_s + g_D \langle \bar{d}d \rangle_0$ is the coupling constant and $\Pi_K(m_K^2)$ is the loop integral in this channel which is given in Appendix C.

Differentiating the dispersion equation, one gets

$$\langle \bar{q}q \rangle_K \frac{\partial \Pi_K}{\partial m_K} = C_{Ku} Q_{ui} + C_{Ks} Q_{si} - 2 \frac{\langle \bar{u}u + \bar{s}s \rangle_0}{(M_+ - m_+)^2} (\delta_{ui} + \delta_{si}), \quad (6.19)$$

where

$$C_{Ki} = 2 \frac{I_i^*(0)}{M_+ - m_+} + 2 \frac{\langle \bar{u}u + \bar{s}s \rangle_0}{(M_+ - m_+)^2} - \frac{\partial \Pi_K}{\partial M_i}, \quad (6.20)$$

with $M_+ = M_u + M_s$ and $m_+ = m_u + m_s$. The explicit forms of $\partial \Pi_K / \partial m_K$ and $\partial \Pi_K / \partial M_i$ are given in Appendix C, again. It is easy to prove that C_{Ki} , which is responsible for a non-vanishing $\langle \bar{d}d \rangle_{K^\pm}$, vanishes in the chiral limit.

A numerical calculation without recourse to chiral perturbation gives

$$\langle \bar{u}u \rangle_{K^\pm} = 1.79, \quad (6.21)$$

$$\langle \bar{s}s \rangle_{K^\pm} = 1.83,$$

$$\langle \bar{d}d \rangle_{K^\pm} / \langle \bar{u}u + \bar{s}s \rangle_{K^\pm} = 8.3 \times 10^{-3},$$

$$\langle \bar{c}c \rangle_{K^\pm} / \langle \bar{u}u + \bar{s}s \rangle_{K^\pm} = 3.6 \times 10^{-3},$$

$$\langle \frac{\alpha_s}{\pi} G^2 \rangle_{K^\pm} = -211 \text{ MeV}.$$

One can see that the sea-quarks are relatively suppressed in the kaon compared with those in baryons. This is again due to the Nambu-Goldstone

nature of the kaon. The anomalous quark content $(\bar{d}d)_{K^\pm}$ and the charm content are comparable to each other and larger than those in pion by an order of magnitude. The former is due to the smallness of the d -quark mass compared to s -quark and the latter is mainly because the kaon mass is larger than pion i.e., the gluon content is larger in the kaon than in the pion.

Using the quark contents obtained, we get the kaon-Higgs coupling,

$$g_{K\phi} = \sum_q m_q (\bar{q}q)_K \cdot \frac{m_K}{v}, \quad (6.22)$$

$$= 301 \frac{m_K}{v} \text{ MeV}, \quad (6.23)$$

which is much larger than $v g_{\pi\phi}$ because the s -quark has a dominant contribution in this case.

7 Summary and concluding remarks

In this paper, we have calculated quark and gluon contents of light baryons and a few pseudo-scalar mesons, and examined the violation of the OZI-rule in the hadrons: baryons are described by the chiral quark model in which the constituent quark masses are given dynamically by the spontaneous breaking of chiral symmetry. Mesons are constructed dynamically as collective excitations described in the model.

As for the light sea-quarks (u, d, s) in the hadrons at low energy scales ($\sim 1 \text{ GeV}$), the main source of non-perturbative OZI-violating process is identified with the determinant 6-fermion interaction due to the axial anomaly, which gives rise to the flavor mixing. The strength of it was determined in the meson sector by fitting the observables of the η - η' system where the anomaly plays a crucial role.

By using the same strength of the vertex, we calculated the anomalous $\bar{q}q$ contents of the hadrons, such as the s -quark content of the nucleon and pion, the u -quark content of Ω , the d -quark content of $K^\pm$ and so on. These

sea contents are closely related to the magnitude of the explicit breaking of chiral symmetry in the hadron system as given by $\bar{\Sigma}_\pi N$.

The application of the simple 't Hooft vertex that was derived for light quarks is not valid for the heavy quarks such as charm. However, by using a sum rule for the hadron masses (the trace anomaly) and the heavy-quark mass expansion, we could extract the heavy-quark contents at the scale of order 1 GeV rather unambiguously. Furthermore, we gave a realistic estimate of the intrinsic charm and gluon contents of hadrons and the Higgs-hadron coupling constants. We have also found that the anomalous quark contents in pseudo-scalar mesons are much smaller than those in baryons; the chiral symmetry and its dynamical breaking are responsible for this fact.

The magnitude of the heavy quark contents of the light hadrons which we have obtained may constitute a basis for a description of the heavy-quark production in terms of the so-called diffractive dissociation[4]. In fact the magnitude of the c -quark in nucleon (pion) was found to be 0.5% (0.03 %) in our estimate which is consistent with the mechanism. We have also shown that the sextet-quark in nucleon, if it exists, is not so much suppressed and of order 0.01%.

The quark contents in the scalar channel $\bar{q}q$ are rather special among those in other channels in the sense that chiral symmetry and a sum rule for hadron masses (the trace anomaly) work quite effectively to extract their values. In other channels, one will have to turn to a model to extract the magnitudes of the quark content, hence the results would become rather model-dependent. Nevertheless, it will be interesting to calculate the quark and gluon contents in the various channels using explicit models of baryons and compare them with our results: the calculation with the MIT bag [44] seems to overestimate the charm content of the nucleon because a large $\alpha_s \sim 2.2$ was used, as we mentioned. Recently it has been found that the heavy quarks could be incorporated in the Skyrme model in the bound-state approach[53]. Such a progress of the model would be most effective in a calculation of the heavy quark content of baryons.

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Appendix A

In this appendix, we give an explicit form of Q_{ji} which turns out to be the scalar meson propagator.

By differentiating the gap equation (2.4) with respect to m_i ($i = u, d, s$), one gets three linear equations for Q_{ji} ($j = u, d, s$) with i fixed. The three equations turn out to be nicely summarised in a matrix equation

$$\mathbf{A} \cdot \mathbf{Q} = \mathbf{I}, \quad (\text{A.1})$$

where

$$(\mathbf{A})_{ji} = \begin{pmatrix} 1 + 2g_s R_u & 2g_D (\bar{s}s)_0 R_d & 2g_D (\bar{d}d)_0 R_s \\ 2g_D (\bar{s}s)_0 R_u & 1 + 2g_s R_d & 2g_D (\bar{u}u)_0 R_s \\ 2g_D (\bar{d}d)_0 R_u & 2g_D (\bar{u}u)_0 R_d & 1 + 2g_s R_s \end{pmatrix}_{ji} \quad (\text{A.2})$$

Here

$$R_i \equiv \frac{d\langle \bar{q}q \rangle_0}{dM_i}, \quad (\text{A.3})$$

$$= -\frac{N_c}{\pi^2} \int^A p^2 dp \frac{p^2}{E_p^3}, \quad (\text{A.4})$$

$$= -\frac{N_c \Lambda^2}{2\pi^2} \left[3\sqrt{1+x^2} - \frac{2}{\sqrt{1+x^2}} - 3x^2 \ln\left(\frac{1+\sqrt{1+x^2}}{x}\right) \right], \quad (\text{A.5})$$

with $E_p = \sqrt{M_i^2 + p^2}$, $x = M_i/\Lambda$ and $\mathbf{I}$ being the identity is originated from $dm_j/dm_i = \delta_{ji}$.

A notable point is that

$$R_i = I_i'(0), \quad (\text{A.6})$$

where

$$I_i'(k^2) = -4N_c i \int \frac{d^4 p}{(2\pi)^4} \frac{p_+ \cdot p_- + M_i^2}{(p_+^2 - M_i^2)(p_-^2 - M_i^2)}, \quad (\text{A.7})$$

with $p_{\pm} = p \pm k/2$ being the loop integral of the i -th quark in the scalar channel. Thus one gets

$$Q_{ji} = [1 + \mathbf{V}_\sigma \cdot \mathbf{\Pi}_\sigma(0)]_{ji}^{-1}. \quad (\text{A.8})$$

[29] W. A. Bardeen, Nucl. Phys. B (Proc. Suppl.) **7A** (1989) 149.

[30] A. Manohar and H. Georgi, Nucl. Phys. **B234** (1984) 189

[31] T. Goldman and R. W. Haymaker, Phys. Rev. **D24** (1981) 724,743.

[32] M. K. Volkov, Ann. Phys. (N.Y.) **157** (1984) 282; A. Dhar, R. Shankar and S. R. Wadia, Phys. Rev. **D31** (1985) 3256; D. Ebert and H. Reinhardt, Nucl. Phys. **B271**(1986) 188.

[33] T. Hatsuda and T. Kunihiro, Prog. Theor. Phys. **74** (1985) 765.

[34] T. Kunihiro, the lectures given in the summer school of Genshikaku-Sausha, held Iiyama 31 July - 5 August 1988, Genshikaku Kenkyu, **33**(1989) 19.

[35] T. Hatsuda and T. Kunihiro, Phys. Lett. **198B** (1987) 126.

[36] G. 't Hooft, Phys. Rep. **142** (1986) 357; M. A. Shifman, A. I. Vainshtein and V. Z. Zakharov, Nucl. Phys. **B163** (1980) 46.

[37] J. Gasser and H. Leutwyler, Phys. Rep. **87** (1982) 77.

[38] M. Takizawa, K. Tsushima, Y. Kohyama and K. Kubodera, Nucl. Phys. **A507** (1990) 611.

[39] T. Hatsuda, Phys. Rev. Lett. **65** (1990) 543.

[40] T. Hatsuda and T. Kunihiro, Prog. Theor. Phys. Suppl. **91** (1987) 284.

[41] H. Pagels, Phys. Rep. **16** (1975) 219.

[42] E. V. Shuryak and J. L. Rosner, Phys. Lett. **218B** (1989) 72.

[43] N. N. Scoccola, Univ. Regensburg Preprint TPR-89-39 (1989).

[44] J. F. Donoghue and E. Golowich, Phys. Rev. **D15** (1977) 3421.

[45] K. Kang and A. R. White, Phys. Rev. **D42** (1990) 835.

[46] W. J. Marciano, Phys. Rev. **D21** (1980) 2425.

[47] K. Fukazawa, T. Muta, J. Saito, I. Watanabe, M. Yonezawa and M. Inoue, Hiroshima Univ. Preprint, HUPD-8917 (1990).

[48] L. I. Reinders, H. Rubinstein and S. Yazaki, Phys. Rep. **87** (1982) 77.

[49] G. E. Brown, K. Kubodera and M. Rho, Phys. Lett. **192B** (1987) 273.

[50] S. Dawson, J. F. Gunion and H. F. Haber, Phys. Rev. **D41** (1990) 2844. K. Hikasa, *Higgs Physics* in Proceedings of the 3rd Meeting on Physics at TeV Energy Scale, (ed. K. Hidaka and C. S. Lim) KEK report 90-9 (1990) p.83.

[51] G. A. Christos, Phys. Rep. **116** (1984) 251.

[52] V. A. Novikov and M. A. Shifman, Z. Phys. **C8** (1981) 43; M. B. Voloshin, Sov. J. Nucl. Phys. **44** (1986) 478; S. Raby and G. B. West, Phys. Rev. **D38** (1988) 3488.

[53] M. Rho, D. O. Riska and N. N. Scoccola, Preprint TPR-90-41 (1990).

Fig. 1 : A schematic figure representing the quark content in baryons. The blob shows the interaction of the external operator $\bar{q}q$ with the constituent quarks. (For light quarks, it is saturated by the propagator of the scalar-meson at zero momentum.) The wavy lines show the short and long range interactions between constituent quarks which affect the quark content.