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TOTAL WIDTHS OF LIGHT HADRONS : AN INCLUSIVE VIEW

A.N. Mitra and Anju Sharma

Department of Physics , University of Delhi, Delhi - 110 007, INDIA

ABSTRACT

The total width of a $q\bar{q}$ hadron is modelled by the rate at which it dissociates into a pair of quasi-free quarks brought about by 'soft' and 'hard' gluon-exchanges, under the ansatz of a short time-zone for the existence of such a state. A crucial ingredient in this regard is the quark's mass function for which a non-perturbative formula is obtained via dynamical breaking of the chiral symmetry of an input vector-exchange-like 4-fermion Lagrangian, facilitated by the knowledge that for a chirally invariant Lagrangian this quantity is also equal to the pion-quark vertex function in the chiral limit ($M_\pi \rightarrow 0$). The mass function so employed is based on a Bethe Salpeter (BS) model with a vector-exchange like kernel (chirally invariant) which is precalibrated to the spectroscopy of $q\bar{q}$ and qqq hadrons, and also predicts a value of $\langle q\bar{q} \rangle_0$ fully consistent with QCD sum rules. The predicted widths show a good overlap with available data from $L > 1$ to $J = 5$, while limitations of phase space resist applicability of the said mechanism due to the onset of hadronic selection rules for certain $L \leq 1$ states.

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1. INTRODUCTION

The total width of a hadron represents one of its fundamental properties which is readily accessible to experiment [1], yet a theoretical understanding of this quantity has, since the inception of the quark model, relied almost exclusively on an "outside-in" strategy to saturate its value through as many observable branching modes as possible. Many authors in the past have attempted to test their models of strong interaction through exhaustive calculations of this kind, a famous example being the FKR model [2]. On the other hand, the intrinsically simpler nature of this quantity (as the

imaginary part of the hadron mass) than such indirect estimations should suggest, perhaps calls for a simpler, "inside-out" point of view, wherein its value is made to depend directly on the distribution of its quark constituents into which it must first break up before the latter should eventually hadronise. Unfortunately the simple dissociation of a hadron into a pair of "free" quarks would not make literal sense because of the obvious confinement problem which in simpler language is tantamount to an infinitely rising potential giving rise to bound states corresponding to real eigenvalues only ($\text{Im. } E = 0$). This is also consistent with the $N_c \rightarrow \infty$ QCD scenario in which limit the total width must vanish. Of course in practice the modest value $N_c = 3$ considerably dilutes the rigidity of this rule, so that alternative strategies for the evaluation of total widths may not be unwarranted.

Is it possible to evolve a workable strategy by means of which an 'acceptable' mechanism for hadronic total widths, directly in terms of their quark / gluon constituents, may be articulated in a quantitative fashion? To that end it is useful to look for some clues from situations where similar mechanisms are believed to be in operation already. Two such situations which exhibit considerable similarity in terms of their actual modes of operation are :

- (i) the break-up of a proton into its partonic constituents due to a hard electromagnetic "kick" characteristic of deep inelastic scattering (DIS) (Fig. 1 a) ;
- (ii) Break-up of heavy quarkonia ($Q\bar{Q} = c\bar{c}, b\bar{b}$) into lighter flavours through the so-called "hidden strong interaction" wherein these two flavour species are connected (internally) by two / three gluon lines for $C = \pm 1$ states respectively (figs. 1 b, 1 c).

In either of these two cases the recombination (reconversion) probability of the (intermediate) quark / gluons into hadrons is effectively unity when one sums over all channels, while the effects of phase space, hadronic selection rules, final state interactions etc. become largely redundant. Thus the 'second stage' process of recombination into hadrons (with probability unity) can be effectively dropped, and the dynamics thus reduced to finding only the 'first stage' probability of dissociation of the hadron into its (i) quark constituents (caused by the e.m. 'kick') in the first case,

and (ii) into its gluonic constituents (albeit in a spontaneous manner) in the second case.

Now if one rushes blindly to apply a similar mechanism to the decay of light hadrons, one immediately encounters serious difficulties. For one thing, the confined nature of the quarks inside the hadron would resist its dissociation into its partonic constituents in the absence of the e.m. kick which characterizes the DIS process (Fig. 1 a). Nor does the gluonic component ($2g / 3g$) have the slightest chance against the dominant $q\bar{q}$ component in the quark / gluon Fock space of a light hadron for dissociation into a similar species, unlike the OZI-forbidden heavy quarkonium ($Q\bar{Q}$) situation which permits the transition $Q \rightarrow q$ only via the gluonic "bridge" (Fig. 1 b, 1 c). Rather, the diagrams for transition of a light quarkonium ($q\bar{q}$) into similar hadronic species would look more like those in Fig. 2 (a,b,c...) in which the two parent quarks ($q, \bar{q}$) are very much a part of the final hadronic scenario, unlike the case of heavy flavoured hadrons wherein the OZI selection rule forces the Q and $\bar{Q}$ to "close in" on themselves as quark loops (Fig. 1 b, 1 c), leaving only the gluons ($2g / 3g$) to 'communicate' with the lighter species through, say, the hooking on of $q\bar{q}$ pairs at the end of each gluon line (as shown).

Taking account of the constraints exhibited in the foregoing, we envisage the following possible mechanism for estimating the total widths of the light hadrons, in a two-step fashion. We assume that there exists a small time zone schematically indicated by a pair of vertical dotted lines in each of Figs. 2 (a,b,c), during which the quark constituents ($q, \bar{q}$) of the hadron stay in dissociated, yet highly confined, state after which they hadronize a la various exclusive modes as are illustrated in Figs. 2 (a,b,c,...). To get a more inclusive view of the entire process, obtained by summing over the individual exclusive modes, we again assume (as in the DIS or the decay of ($Q\bar{Q}$) hadrons) that the second stage hadronization occurs with probability one, so that it is enough to concentrate on the first stage zone (indicated by the dotted lines in fig. 2) only. This is depicted separately in Figs. 3 (a,b) in which the confined quarks (mass m_q) inside the parent hadron get converted to a quasi free, yet

confined, state (mass m_f) through the exchange of a soft gluon (non-perturbative) and hard gluon (perturbative) respectively. It is the latter state of the quark (mass m_f) that is deemed to participate in the eventual hadronization with probability one, through a summation over all possible exclusive modes a la Figs. 2 (a,b,c) etc. Admittedly such a mechanism for light flavour transitions is not as "clean" as the cases of DIS or heavy flavour transitions noted above, since it depends crucially on the ansatz of a (small) time-zone during which the quark pair has a quasi-free existence before final hadronization, yet the mechanism has an intuitive appeal which is persuasive enough to warrant a closer scrutiny in terms of direct comparison with data. To that end the detailed dynamics of the process needs several ingredients - the vertex function $\Gamma(q)$ at the parent hadron-quark vertex, the momentum dependence of the mass function $m(p^2)$ which gives a formal meaning to the quantities m_q and m_f noted above, and of course the availability of the phase space which is determined by the mass (M) of the parent hadron.

At this stage it is useful to give a comparison of Figs. 3 (a, b) for spontaneous hadronization with the DIS process Fig. 1a : The mass m_f of the quasi-free quarks in the former corresponds to the much smaller 'current' mass (m_c) of the partons in the latter (brought about by the e.m. kick), but is now expected to have a substantially higher value (due to confinement), roughly intermediate between the constituent mass (m_q) and the current mass (m_c). The precise value of m_f will ofcourse be determined (see Eq.(22) below) dynamically by the running mass $m(p^2)$ with due regard to formal energy-momentum conservation in Figs. 3 (a, b). The momentum dependence of $m(p^2)$ is in turn governed by the Nambu-Jona-Lasino (NJL) mechanism [3] for dynamical breaking of chiral symmetry (DB χ S), or its Schwinger-Dyson Equation (SDE) counterpart in conventional field theory [4,5], where the emphasis on the non-perturbative aspect of the problem necessitates a conscious departure from the perturbative QCD regime which corresponds to a mass formula of the form [6,7]

$$m_{\text{QCD}}(-p^2) = 4 \pi \alpha_s (-p^2) \langle q\bar{q} \rangle_0 / 3 (-p^2) \quad (1)$$

in favour of an alternative structure based on non-perturbative premises. We address this issue in the next section (2) in the

context of a concrete dynamical model based on an interplay of the Bethe Salpeter Equation (BSE) for hadron spectra [8,9] and the solution of the Schwinger-Dyson Equation (SDE) for the quark mass function $m(p^2)$ which satisfies DB χ S [10]. However before going into these details, it is first necessary to check more explicitly on the compatibility or otherwise of the proposed mechanism with the $N_c \rightarrow \infty$ scenario of QCD.

$N_c = \text{infinity limit} :$

Fig. 3(a,b) represents the generic diagram for our proposed hadronisation mechanism via suitable quark loops built around it, as illustrated by Figs. 2 (a,b,c,...) which consist of a sequence of single quark loops of n sides ($n = 2, 3, 4, 5, \dots$) obtained by inserting $(n-1)$ $q\bar{q}$ pairs. The sum of these diagrams which symbolize a mechanism of first stage hadronization should be illustrative enough of the N_c -dependence of the total hadronic amplitude in the $N_c \rightarrow \infty$ limit even without going in for more involved diagrams representing second and higher stages of hadronization. The N_c -dependence of these individual loop diagrams in turn may be worked out by noting that at each $Hq\bar{q}$ vertex the factor $\delta_{\alpha\beta} / \sqrt{N_c}$ is operative [11], where $\alpha, \beta (=1,2,3)$ are the color indices of the quarks involved. This is a precise $3 \rightarrow N_c$ generalization of the color dependence of such $Hq\bar{q}$ vertices which had been previously employed [12] for evaluating the corresponding loop integrals for several exclusive hadronic decays within a BS model for $Hq\bar{q}$ vertex functions [12], and since the same model is also envisaged in the present context (see below) there is no danger of wrong color counting. The pion-quark vertex function is no exception to this $\sim N_c^{-1/2}$ rule [11,12], an obvious assertion necessitated by the dominant role of pionic modes in hadronic decays [1].

For a crude estimate of the total decay amplitude, the remaining dynamical factor at the main $Hq\bar{q}$ vertex and the subsequent ones may be notionally denoted by f and g respectively. Thus the successive loops 2(a,b,c,...) give rise to the sequence of amplitudes $f g^2 / \sqrt{N_c}$; $f g^3 / N_c$; $f g^n / N_c^{(n-1)/2}$ whose sum (as a geometric series) yields the value

$$A_H = f g^2 / [\sqrt{N_C} - g] \quad (2)$$

so that the total hadronic amplitude A_H does indeed behave asymptotically as $1 / \sqrt{N_C}$. This $N_C \rightarrow \infty$ result is rather insensitive to local variations in the weighting procedure for the successive loops diagrams ; e.g. weighting by, say, $(n!)^{-1}$ will not alter this qualitative conclusion. On the other hand the generic diagram Fig. 3 (a,b) which is inclusive in character, would prima facie seem to go against this conclusion since the corresponding amplitude has the color dependence $f \bar{u}_\alpha u_\alpha / \sqrt{N_C}$, which in turn would give rise to a total width proportional to $|f|^2$, apparently independent of N_C (?), were it not for the fact that the vertex function $|f|$ is also of an identical nature to the mass function $m(p^2)$ which in the perturbative QCD regime already goes down like N_C^{-1} , (see eq.(1)), since the factor α_s is itself proportional to $1/N_C$ according to the QCD wisdom of t'Hooft and Witten [11]. On the other hand, the non-perturbative formula for $m(p^2)$ that we are going to employ gives it a much sharper fall-off, viz., $\exp(-\sqrt{N_C} A)$ with N_C (see Eq.(7) below), so that our "inclusive strategy", vide Figs.(3a,3b), stays well within the $N_C \rightarrow \infty$ rule.

After this vital color check, we turn to a quantitative description of Figs. 3(a,b) to prepare the ground for an experimental test of our ansatz that this diagram with two outgoing quarks of appropriate mass (m_f) governed by 4-momentum conservation (see below), already stands for the entire hadronization process, without explicit evaluation of the individual hadronic modes. This picture of hadronization of light quarkonia, though partly motivated by the OZI-forbidden

heavy ($Q\bar{Q}$) $\rightarrow$ light ($q\bar{q}$) transitions

Figs. 1(b,c), and partly by the DIS process, Fig. (1a), is admittedly on less firm grounds (because of the additional ansatz of a small time zone) ^{for quasi-free existence} but appears plausible enough to stand further scrutiny in terms of actual comparison with data.

A spectroscopy-related approach based on DBxS :

To sum up, the Fock space responsible for the eventual hadronic transition is assumed to be dominated by a pair of quasi-free quarks

(mass m_f each), in comparison to other components of Fock-space, such as $2g / 3g$ and / or $(q\bar{q})^2$ which, though verifiable in principle, are expected to have much less importance in the (present) context of light flavour transitions than, e.g., for heavy flavour transitions. The diagrams (3a) and (3b) stand for the effects of soft and hard gluon exchanges respectively which convert the mass m_q "inside" the hadron to the quasi-free mass m_f "outside". The latter is expected to have a value somewhat in between the constituent mass m_q and the current mass m_c . The change from m_q to m_f , in turn is rather naturally accounted for in terms of a momentum-dependent mass function $m(p^2)$ of the quark, brought about through a non-perturbative mechanism based on dynamical breaking of chiral symmetry (DB χ S). The other ⁰¹important ingredient for the evaluation of the transition amplitudes (3a,3b), and the determination of their relative importance, is the vertex function $\Gamma(q)$ sitting at the parent hadron-quark junction. Both these quantities have been obtained in a unified fashion through an interlinked BSE-cum-SDE mechanism wherein the BSE serves as the basic vehicle for addressing all the flavour sectors of hadron spectroscopy [8,13] in terms of a vector-exchange-like confining kernel (which automatically obeys chiral symmetry). With the chirally symmetric $q\bar{q}$ interaction so defined, which goes beyond the contact type NJL interaction [3], the dynamical determination of the mass function $m(p^2)$ as a non-trivial solution of the SDE no longer needs a cut-off, since the confinement parameter now automatically plays the role of the latter. Alternatively, the BSE, when specialized to the pion in the chiral limit ($M_\pi = 0$), is known [4,5] to reduce to the SDE for the mass function $m(p^2)$, thus providing an independent handle on this crucial quantity via the spectroscopy sector. This property was recently exploited [10] for a non-perturbative construction of $m(p^2)$, see Eq. (17) below .

All these results on $\Gamma(q)$ and $m(p^2)$ which are sought to bring out the vital role of the spectroscopy sector in determining these basic functions, have been adequately described in some recent [8-10] and not so recent [12,13] publications. However since the main thrust of this spectroscopy-oriented approach does not seem to have received enough attention in the literature, it may be useful to first

summarise the basic philosophy as well as the main findings, in the next section (sec. 2) before describing the evaluation (in sec. 3) of the total widths through an interplay of the diagrams (3a) and (3b) representing the soft and hard gluon contributions respectively. In sec. 4 the numerical results are presented after taking due cognizance of the intrinsic limitations of this model for comparison with the actual data.

2. A BSE-cum-SDE Program for $\Gamma(q)$ and $m(p^2)$:

The basic philosophy of this program, as outlined in [10] in relation to the contemporary literature, is to make simultaneous contact with spectral data on ^{the} one hand and transition amplitudes on the other, using the vehicle of the BSE whose kernel consists of an effective confining term (vector-exchange-like) playing the role of a soft gluon exchange (taken non-perturbatively), and a one-gluon-exchange term (taken perturbatively). Taking account of the $O(3)$ -like nature of the spectra [11], it had been found necessary at the outset, to consider the BSE in a two-tier fashion : the first stage reduction to a 3D form [16] is appropriate for making contact with the spectra, while the reconstructed 4D wave function, or equivalently the vertex function $\Gamma(q)$ facilitates the evaluation of transition amplitudes via appropriate quark loops [12]. A consistent mathematical strategy for performing this twin function was to require a 3D support to the BS kernel, albeit in a covariant fashion which was achieved somewhat later [9] and provided an exact interconnection between the 3D and the 4D forms of the BSE [9]. Thus the structure of the vertex function $\Gamma(q)$ could be 'determined spectroscopically' in terms of the input parameters of the BS kernel expressed as a covariant 3D 'potential' which has the following value for equal mass kinematics [13,8] :

$$V(r) = M_{>} \omega_0^2 \alpha_s(M_{>}^2) \left[r^2 / (1 + A_0 M_{>}^2 r^2)^{1/2} - C_0 / \omega_0^2 \right] \quad (3)$$

$M_{>} = \max. (M, 2 m_q)$, m_q being the constituent quark mass ; $\alpha_s(M_{>}^2)$ is the strong coupling constant ; ω_0 is a reduced spring constant having the dimensions of a string tension ; C_0 a dimensionless constant simulating the effects of zero point energy, and A_0 is a small constant which acts as an interpolating device between the harmonic

form ($\sim r^2$) for small r and a linear form ($\sim r$) for large r . This potential, though constructed empirically, has all the "good" properties of a standard confining potential, viz. it behaves linearly for large r with a proportionality factor ($\omega_0^2 \alpha_s / \sqrt{A_0}$) independent of flavour mass (except for logarithmic variations due to α_s). For small r on the other hand it shows a quadratic behaviour which is again consistent with theoretical expectations based on a slowly varying background gluon field [14]. The parameters ω_0 , C_0 and A_0 were determined [13,8] from spectroscopy [1] as

$$\omega_0 = 158 \text{ MeV} ; \quad C_0 = 0.27 ; \quad A_0 = 0.028 \quad (4)$$

while the strong coupling constant was taken as [8]

$$\alpha_s = 6 \pi / (33 - 2f) \log (M_{>} / \Lambda) ; \quad M_{>} = \max (M, 2m_q) ; \quad \Lambda = 200 \text{ MeV} \quad (5)$$

For each flavour sector the constituent mass m_q was determined from spectroscopy [13,8], the value of m_q for u - d -quarks being found as 265-270 MeV [8,13], but the latter was subsequently determined in a dynamical fashion through a self-consistent solution of the SDE [10] to yield a value of 300 MeV corresponding to the value of the mass function $m(p^2)$ at $p = 0$. Thus it was no longer necessary to regard it as a free parameter, in the spirit of the Nambu-Jonalasino (NJL) model [3] based on the principle of chiral invariance at the input Lagrangian level and its breaking through a non trivial solution of the Schwinger-Dyson equation (SDE) (the gap equation). In the present case the NJL spirit has been implemented by extending the contact-type interaction to gluon-exchange-like 4-fermion interactions which also share the chiral invariance property because of their vector-exchange character) [4,5,7]. The spatially extended nature of the latter moreover brings out more clearly another important property that was obscured from view in the contact-type NJL model [3], viz., the equality of the mass function with the pion-quark vertex function in the chiral limit, as a direct consequence of the vector-exchange type (chirally invariant) interaction [4] via Ward-Takahashi (WT) identities [5]. This involves a non-trivial interplay of SDE and Bethe-Salpeter equation (BSE) and provides a concrete handle on $m(p^2)$ through the BS sector of the dynamics. Our use [10] of this last property is clearly justified within a BS model characterized by a vector-exchange type kernel

[8,13] which automatically enjoys this built-in chiral property.

Structure of the Vertex Function $\Gamma(q)$:

To proceed further, the interconnection between the 3D and 4D forms of the BS wave function brought about by the ansatz of a 3D support to the BS kernel [9], gives rise to a structure of the form $D(\hat{q}) \times \phi(\hat{q})$, for the hadron-quark vertex function, where $D(\hat{q})$ and $\phi(\hat{q})$ are respectively the denominator function and wave function of the 3D BSE which resembles a Schroedinger equation (albeit in a covariant form). For a harmonic kernel which follows from Eq. (3) for $A_0 = 0$ (which is adequate for ud flavours), the complete structure of the vertex function for different types of L-excited hadrons is expressible in terms of the input parameters ω_0 , C_0 and m_q in the following form which does not yet take account of the o.g.e. effect (see sec. 3 below) as follows :

$$\Gamma_H(\hat{q}) = N_H D(\hat{q}) \frac{\phi(\hat{q})}{(2\pi)^{5/2}} \Gamma_i ; \phi(\hat{q}) = \exp(-\hat{q}^2 / 2 \beta_H^2) , H \neq \text{pion} \quad (6)$$

$$\beta_H^4 \equiv \beta^4 = \frac{1}{2} M_{>}^2 \omega_0^2 \alpha_s(M_{>}^2) / \alpha_N (1 - 2 \alpha_s C_0) ; \alpha_N \sim 0(1) \quad (7)$$

$$D(\hat{q}) = 4 \hat{\omega} (\hat{\omega}^2 - M^2/4) ; \hat{\omega}^2 = m_q^2 + \hat{q}^2 \quad (8)$$

$$2q = p_1 - p_2 ; P = p_1 + p_2 = k_1 + k_2 ; \hat{q}_\mu = q_\mu - q \cdot P P^{-2} P_\mu \quad (9)$$

Γ_i is a Dirac-matrix which has the following tensorial structure for the various physical states of interest, corresponding to $S = 0$ (B_J with $J = L$) and $S = 1$ (V_J , A_J , S_J with $J = L + 1, L, L-1$ respectively)

$$B_J(\pi, b_1, h_1, \pi_2) \rightarrow \gamma_5 B^J(\hat{q}) ; \quad (10)$$

$$B^J(\hat{q}) = (2^L / L!)^{1/2} \hat{q}_{\mu_1} \dots \dots \hat{q}_{\mu_J} B_{(\mu)}^J \beta^{-J}$$

$$V_J(\rho, \rho_3, \omega_3, f_2) \rightarrow i \gamma \cdot \mathcal{E}^J(\hat{q}) ; \quad (11)$$

$$\mathcal{E}_\mu^J(\hat{q}) = (2^L / L!)^{1/2} \hat{q}_{\mu_1} \dots \dots \hat{q}_{\mu_{J-1}} \mathcal{E}_{(\mu)}^J \beta^{-J+1}$$

$$A_J (a_1, f_1 \dots) \rightarrow i\gamma_5 \gamma \cdot \mathfrak{z}^J(\hat{q}) ; \quad (J=1 \text{ only}) \quad (12)$$

$$S_J (f_0, a_0 \dots) \rightarrow i\gamma \cdot \hat{q} B^J(\hat{q}) ; \quad (J=0 \text{ only}) \quad (13)$$

We have indicated a few typical states [1] against each symbol.

$B_{(\mu)}^J$, $\mathfrak{z}_{(\mu)}^J$ are tensors of appropriate rank J satisfying the following normalizations

$$\begin{aligned} \langle B^J(\hat{q})^* B^J(\hat{q}) \rangle &= (2 \hat{q}^2 / \beta^2)^J / (2J+1)!! \equiv f_J \\ \langle \mathfrak{z}_{\mu}^J(\hat{q})^* \mathfrak{z}_{\mu}^J(\hat{q}) \rangle &= f_{J-1} \end{aligned} \quad (14)$$

The BS normalizer in turn is given for equal masses by [9]

$$N_H^{-2} = \int d^3 \hat{q} \phi^2(\hat{q}) 2 \hat{\omega} (4 \hat{\omega}^2 - C_J M^2) T_J(\hat{q}) \quad (15)$$

where $T_J(\hat{q})$ and $C_J(\hat{q})$ are given for the four different cases as under:

B_J	V_J	A_J	S_J
$T_J(\hat{q}) = f_J$	f_{J-1}	$2f_{J-1}$	$(2J+3)f_{J+1}/(J+1)$
$C_J(\hat{q}) = 0$	$[4J/(2J+1)]\hat{q}^2/M^2$	$[4J/(2J+1)]\hat{q}^2/M^2 + 4m_q^2/M^2$	$4 \hat{q}^2/M^2$

(16)

In all these expressions the 'hat' on the internal 4-momentum q_{μ} stands for its perpendicularity to the hadron 4-momentum P_{μ} , a result which reflects the 3D support to the BS kernel [9].

The Non-Perturbative Mass Function $m(p^2)$:

From the structure of the vertex function Γ_H , eqs.(6,10-13) above, it is now possible to infer the structure of the mass function $m(p^2)$ by specializing the former to the case of the pion in the chiral limit ($M_{\pi} = 0$), in accordance with the chiral invariance property of the vector-exchange kernel noted above. This has been described in detail in ref. [10], where it was explicitly shown how in this limit the SDE and the BSE become identical, but the main points may be recapitulated for the sake of completeness : Clearly it is now enough to deduce the mass function from $\Gamma_{\pi}(\hat{q})$, eq. (6), except for the

overall normalization which must be fixed anew. In this regard we continue to adhere to the prescription of [10], viz., we demand that the mass function reduce to the constituent mass m_q in the $p = 0$ limit, i.e., $m(0) = m_q$. Thus the mass function acquires the non-perturbative form (note that for the pion, $M_\pi = 2 m_q$) :

$$m(\hat{p}_\pi^2) = \frac{\hat{\omega}_\pi^3}{m_q^2} \exp(-\hat{p}_\pi^2 / 2 \beta_\pi^2) ; \quad (17)$$

$$\beta_\pi^4 = 2 m_q^2 \omega_0^2 \alpha_s(4 m_q^2) / (1 - 2\alpha_s C_0) ; \quad m_q = m(0) \quad (18)$$

$$\hat{\omega}_\pi^2 = m_q^2 + \hat{p}_\pi^2 ; \quad \hat{p}_\pi^\mu = p_\pi^\mu - p_\pi \cdot \eta \eta^\mu ; \quad \eta_\mu = \frac{P_\mu}{\sqrt{P^2}} . \quad (19)$$

Note that the symbol 'p' for each quark momentum $p_{1,2}$ now equals the relative momentum q inside the hadron (pion) in the limit $P_\mu = 0$. The symbol ' $\hat{}$ ' on $\hat{p}_\pi$ in the structure of $m(\hat{p}_\pi^2)$ reflects its derivation as a solution of the (two-body) BSE where the 'hadron' is now the pion, albeit in the (chiral) limit of $P_\mu = 0$. This is stressed by the subscript ' π ' on the quark 4-momentum p_μ , a feature which does not exist in the corresponding perturbative formula (1). This has a simple interpretation when the mass function $m(p^2)$ is looked upon as a solution of the SDE wherein the 'hat' on 'p' reflects the presence of a second source of color charge, an information carried by the direction η_μ of the (composite) pion 4-momentum P_μ . This feature which generalises [10] a previously known result [15] in heavy-light quarkonia ($Q\bar{q}$) to any relativistic composite, lends a degree of consistency to the interpretation of m_q as the constituent mass necessitating the presence of a second source of color (see ref. [10] for more details).

As a check on the numerical accuracy of the mass function defined above, we find that the quark condensate defined by [10]

$$\langle q\bar{q} \rangle_0 = N_c N_f \frac{-i}{(2\pi)^4} \int d^4 p \text{Tr } S_F'(p) = \frac{3}{2\pi^3} \int d^3 \hat{p} \frac{m(\hat{p}^2)}{[m^2(\hat{p}^2) + \hat{p}^2]^{1/2}} \quad (20)$$

yields a parameterfree determination $2 \times (266 \text{ MeV})^3$, in fairly good accord with the QCD sum rule value [16] of $\langle u\bar{u} \rangle = (225 \text{ MeV})^3$, when it is remembered that the quantity (20) stands for $\langle u\bar{u} \rangle + \langle d\bar{d} \rangle$, while the QCD-SR value is only for $\langle u\bar{u} \rangle$.

At this stage a word is in order regarding the difference in the value of $\langle q\bar{q} \rangle$ obtained above from that of ref. [10]. This is traceable to the difference in parametrization for β_π^4 , eq.(18), from that of ref. [10] in so far as the hadron mass M_π is now replaced by $2 m_q$. This is formally necessary before going to the chiral limit $M_\pi = 0$, and is also in accordance with the correspondence $2 m_q$ (CIA) $\leftrightarrow$ M (NPA) between the CIA model [9] and the NPA model [8]. For all other hadrons the correspondence $2 m_q \leftrightarrow M$ makes little difference anyway [17]. This modification for the pion case was unfortunately missed in the derivation of the condensate $\langle u\bar{u} \rangle$ which thus gave a lower value (yet consistent with QCD-SR).

Before closing this section we check on the N_c -dependence of $m(\hat{p}^2)$ in eq.(17) which comes from the proportionality of β_π^4 to α_s , see eq.(18). Since α_s , in the t'Hooft-Witten picture [11], is itself proportional to $1/N_c$, this gives β_π^2 proportional to $1/\sqrt{N_c}$, a result which when substituted in the exponential factor on the RHS of eq.(17), yields the proportionality factor

$$m(\hat{p}^2) \sim \exp (- \sqrt{N_c} A) ; A = \text{independent of } N_c \quad (21)$$

This N_c -variation of the non-perturbative $m(\hat{p}^2)$ is clearly stronger than that of the perturbative formula [6,7] given by eq.(1), which suggests merely $m(p^2) \sim 1/N_c$. Since the same N_c -dependence also holds for the hadron-quark vertex function Γ_H , eq.(6), this result checks, without further ado, the expected vanishing of the inclusive total width of a hadron in the infinite N_c limit, following the argument after Eq. (2) . We now come to the evaluation of the total width in an inclusive fashion for the actual case of $N_c = 3$.

3. TOTAL HADRON WIDTH via $H \rightarrow q\bar{q}$ (Figs. (3a, 3b) :

The non-perturbative mass formula (18) can now be used to determine the mass m_f of the pair of quasi-free quarks in Fig. 3(a,b) in accordance with 4-momentum conservation which in the c.m. system of the hadron (mass M) reads as :

$$Q^2 + m_f^2 = M^2/4 \quad (22)$$

where $\pm \vec{Q}$ are the c.m. momenta of the quarks on dissociation. To

determine m_f from eq. (18), we employ the simplest isotropy ansatz on the 4-direction η_μ of the pion in eq. (19), so as to yield $\hat{p}_\pi^2 = -(3/4) m_f^2$, which gives from eq. (18), the estimate $m_f = 176 \text{ MeV}$, which is (expectedly) less than the constituent value m_q inside the hadron. The quasi-free value $-(3/4) m_f^2$ for $\hat{p}_\pi^2$ may be contrasted with the result $p^2 = -m_f^2$ usually employed in perturbative calculations [18] of $m(p^2)$ which do not 'see' the second source of color charge.

The calculations have been performed in accordance with figs. 3(a,b) which interfere coherently and are termed soft (NP) and hard (P) gluon-exchange contributions to the $q\bar{q}$ dissociation respectively. There is also a two-gluon emission amplitude, analogously to heavy to light flavour transitions, Figs.1(b,c), which when relevant would add incoherently to Figs.3(a,b), but the former in general have little chance against the latter in the case of light to light transitions, and so have not been considered in this paper.

For the evaluation of the total widths, the logical strategy for inclusion of the two diagrams Fig.(3a,3b) which correspond to the effects of the confining (h.o.) and one-gluon-exchange (o.g.e) terms respectively in the Bethe-Salpeter kernel [8,13], is taken to be the same as that employed earlier for the calculation of the mass (M) spectra of hadrons [8,13], since the mass (M) and the width (Γ_t) are essentially on the same footing, being the real and imaginary parts of the (complex) energy eigenvalue of a hadron. To recapitulate the strategy in respect of mass (M) spectra [8], the confining (h.o.) term is treated exactly, and the o.g.e. term perturbatively. This procedure may be justified on the ground that the "pole" structure is entirely due to the former, while the o.g.e term merely causes a modest shift in the position of M. The relative signs of these two terms are also correctly reflected in this procedure, since the intrinsic weakness of the o.g.e. term prevents it from generating a pole structure on its own, so that its higher order effects can make at most marginal contributions to M. This practical strategy, which makes sound physical sense, has the added virtue of providing a fully analytical structure without sacrificing numerical accuracy, with the extra confidence gained through the prior agreement of the calculated M-values with the data [8,13].

We shall therefore follow an identical strategy in respect of the total width (Γ_t) calculations which are logically the imaginary counterparts of the mass (M) positions. Thus the soft gluon contribution, Fig.(3a), will be treated exactly, i.e. through an infinite number of iterations of the exchanged gluon line, which amounts to evaluating the entire Hadron-quark vertex function via the BSE with the confining part, eq. (3), of the kernel [13, 8]. For the evaluation of the o.g.e. effect, Fig.(3b), this may be done in the standard perturbative quantum mechanical manner by considering the exchanged gluon line only once (hard gluon effect), while the $Hq\bar{q}$ vertex in the same is taken as in Eq.(6), i.e., the one determined by the confining part of the kernel. The central ingredient is thus the $Hq\bar{q}$ vertex function which for an L-excited hadron is listed in Eqs. (6,10-13)

In accordance with the foregoing discussion on the relative strengths of the contributions from figs.(3a, 3b), the total amplitude for a $H \rightarrow q\bar{q}$ transition is now expressible as :

$$(2\pi)^{-5/2} N_H \bar{u}(k_1) (A_{NP} + A_P) v(-k_2) ; \quad (23)$$

where

$$A_{NP} = \frac{1}{i} D(Q) \phi(Q) \Gamma_i (2\pi)^4 ; \quad 2Q = k_1 - k_2 ; \quad Q \cdot P = 0 \quad (24)$$

$$A_P = (-\frac{4}{3} g_s^2) \int d^4q \gamma_\mu S_F(p_1) \Gamma_i \frac{D(\hat{q})\phi(\hat{q})}{i} S_F(-p_2) \gamma_\mu (-i) q_1^{-2} \quad (25)$$

are the non-perturbative (soft) and perturbative (hard) gluon contributions respectively. Further, $g_s^2 = 4\pi\alpha_s(M^2)$, and Γ_i is listed in (10-13) above for the different cases considered ; $q_1 = q - Q$ is the exchanged momentum in Fig.(3b) out of which $\hat{q}_{1\mu}$ is its 3-D component orthogonal to P_μ , so that since $P \cdot Q = 0$,

$$q_1^2 = \hat{q}_1^2 - M^2 \sigma^2 ; \quad \sigma = q \cdot P / P^2 ; \quad \hat{q}_{1\mu} = \hat{q}_\mu - Q_\mu \quad (26)$$

The apparent simplicity of the non-perturbative amplitude A_{NP} , Eq.(24), compared to the perturbative contribution A_P is due to the following circumstance : If A_{NP} were written in an analogous form to A_P , Eq.(25), the resulting integral would immediately simplify (after

σ - integration) to the RHS of a 3D BSE, viz., $D \phi = \int K \phi$ [9], thus justifying its replacement by the LHS, as indicated by eq. (24). It may be checked that the function $D(\hat{Q})$ is proportional to $m_q^2 - m_f^2$, by virtue of Eq. (22). The evaluation of A_P is considerably simplified by employing Gordon Reduction [19,8] for the two γ_μ matrices in Eq. (25), facilitated by the appearance of the quasi-free quark spinors in Eq. (23). As a result, Eq. (25) is reexpressible as

$$A_P = \left(-\frac{4}{3} g_s^2\right) \int d^3 \hat{q} M d\sigma D \phi [GR] / (\Delta_1 \Delta_2 q_1^2) \quad (27)$$

$$[GR] = [(2k_1 + q_1)_\mu + i \sigma_{\mu\nu} q_{1\nu}] \Gamma_\nu [(2k_2 - q_1)_\mu - i \sigma_{\mu\lambda} q_{1\lambda}] \quad (28)$$

where we have changed the integration variables as $d^4 q = d^3 \hat{q} M d\sigma$ and used the relations $p_1 = k_1 + q_1$ and $p_2 = k_2 - q_1$. Further,

$$\Delta_{1,2} = m_q^2 + p_{1,2}^2 = \hat{\omega}^2 - M^2 (1/2 \pm \sigma)^2 \quad (29)$$

[GR] has a very compact structure for all the cases involved, viz:

$$V_J ; A_J : (-M^2 - 4Q \cdot \hat{q}) (\gamma \cdot \mathcal{X}^J(\hat{q}) [1 ; \gamma_5] + 4 \mathcal{X}^J \cdot Q m_f Q^{-4} (\hat{q} \cdot Q - Q^2)^2 [1 ; -\gamma_5]) \quad (30)$$

$$B_J ; S_{J=0} : (-M^2 - 4Q \cdot \hat{q}) [\gamma_5 B^J(\hat{q}) ; i \gamma \cdot \hat{q} (2/\beta^2)^{1/2}] + q_1^2 [2 \gamma_5 B^J ; 4 i \gamma \cdot \hat{q} (2/\beta^2)^{1/2}] \quad (31)$$

We now integrate A_P w.r.t. $d\sigma$ by contour integration (since the vertex function $D \otimes \phi$ does not involve σ), using the results

$$|\hat{q}_1| D I_0 = (2\omega + |\hat{q}_1|) / [(\omega + |\hat{q}_1|)^2 - M^2/4] \quad (32)$$

$$D I_2 = \hat{q}_1^2 I_0 D - 1 ; I_1 = 0 \quad (33)$$

$$\text{where } 2\pi i I_n \equiv \int (M \sigma)^n M d\sigma / (\Delta_1 \Delta_2 q_1^2) \quad (34)$$

Next the integration over the azimuthal angle ϕ for the 3-vector $\hat{q}_\mu$ is carried out by using the formula

$$\langle \mathcal{X}_\mu^J(\hat{q}) ; B^J(\hat{q}) \rangle_\phi = P_L(\mu) [\mathcal{X}_\mu^J(Q) ; B^J(Q)] \quad (35)$$

where $\cos^{-1} \mu$ is the angle between $\hat{q}_\mu$ and $Q_\mu = \hat{Q}_\mu$. Finally the

integration over $d^3\hat{q} = 2\pi d\mu \hat{q}^2 d\hat{q}$ is carried out numerically (Gauss-Legendre for μ and Gauss-Laguerre for $\hat{q}$) for the different cases involved. As a result all the A_P amplitudes are compactly expressible in terms of two quantities F_L and G_L as

$$A_P (V_J; A_J) = iF_L \gamma_5 \delta^J [1 ; \gamma_5] - iG_L \delta^J \cdot Q m_f Q^{-2} [1 ; -\gamma_5] \quad (36)$$

$$A_P (B_J; S_0) = iF_L \gamma_5 B^J(Q) ; -iF_1 m_f (2 / \beta^2)^{1/2} \quad (37)$$

where

$$F_L ; G_L = (2\pi)^2 (4 g_s^2 / 3) \iint \hat{q}^2 d\hat{q} d\mu \phi P_L(\mu) (\hat{q}/Q)^L 1_0 D * \\ * [(M^2 + 4 Q \hat{q} \mu) ; 4 (\hat{q} \mu - Q)^2] \quad (38)$$

In this notation the total width for V_J states works out as

$$\Gamma(V_J) = N_H^2 (2\pi)^{-6} (Q/2) T_J [(1-C_J) F_{L1}^2 + (G_L^2 - 2F_{L1} G_L) \frac{4J}{2J+1} (m_f/M)^2] \quad (39)$$

where $F_{L1} = F_L - (2\pi)^4 D * \phi$, and T_J, C_J are as defined in eqs. (16), of the text. The other formulae are omitted for brevity.

4. RESULTS AND DISCUSSION

Table 1 summarizes our findings for a fairly representative cross section of non-strange $q\bar{q}$ states which prima facie offer the most natural candidates for testing the veracity of the proposed mechanism using the mass function, Eq. (17), since it almost saturates the constituent mass values of u, d quarks (small current masses). To adapt the formula to other states would require substantial corrections for the strange quark (current) masses, so this question is best relegated to a later stage. On the other hand an uncritical comparison with the data[1] is beset with several formal difficulties associated with the very limitations of the postulated mechanism vis-a-vis the actual situation. Therefore a prior anticipation of cases where the predictions are likely to go 'obviously wrong', as well as those cases where they may be 'not so wrong' should be of some help in judging the "physics" of the mechanism under study. To that end it is necessary to remember that

the observed total widths represent the resultant hadronisation effect of the (quasi-free) quarks, and therefore must be susceptible to hadronic selection rules which the proposed mechanism has made no attempt to address so far. For example, ρ has a large width via the two-pion mode while the ω width is strongly inhibited, below the three-pion threshold. The interpretation of the data is indeed greatly facilitated through a two-way classification in terms of natural versus unnatural parity states ($N^\pm$) on the one hand, and positive versus negative G parity states ($G^\pm$) on the other. Such a classification is also compatible with the spirit of exchange degeneracy [20]. The "best" cases for comparison with the data, viz. those which are least likely to be affected by such hadronic selection rules, are perhaps the N^+G^+ states, while the other pairs are more likely to be susceptible to these selection rules in varying degrees. Their difference from N^+G^+ states should be particularly noticeable for low excitations ($L \leq 1$), but this distinction should rapidly disappear for higher excitations ($L \geq 2$) as more and more channels open up. If the data are viewed in this light, there does appear to be some degree of consistency in the observed pattern. Thus the rather good accord of our results with the data for ρ_3 and ω_3 as well as for π_2 brings out the manner in which the relative effects of the hadronic selection rules have been overcome with increasing L-values. Indeed the general agreement of the calculated widths of V_J states is quite impressive from $N = 2$ to $N = 5$, (for which the data are available). $N = 1$ states occupy an intermediate position between $N = 0$ and $N = 2$ in terms of their susceptibility to the hadronic selection rules. These have been shown more by way of illustrating the "degree of disagreement" (with the predictions) of these cases than with any serious expectation of good agreement since, by virtue of their peculiar configurations-- vis-a-vis the available phase space for hadronization-- their behaviour is bound to be erratic. Yet the (NG) classification does help understand the nature of the discrepancies to some extent. Thus even the $N = 1$ case of $f_2(1270)$ which has an N^+G^+ status, shows good overlap with data, while the discrepancies for b_1 and a_2 are examples of their N^-G^+ and N^+G^- status respectively with low ($N = 1$) excitation. On the other hand, the cases of a_0 and f_0 are perhaps on a different footing, and may not even be amenable to understanding on the basis of a mere (NG)

classification. Their "agreement" with the data is probably an artifact of very limited phase space and not a causal effect of the proposed mechanism. A possible conjecture in this regard is the role of a $(q\bar{q})^2$ configuration in the Fock-expansion, which normally gives a lower order effect than the (main) $q\bar{q}$ term, but may give a significant contribution when the $q\bar{q}$ channel is inhibited by selection rules operating under low phase space conditions.

Encouraged by these results we have also ventured to make some semi-quantitative estimates for the decays of strange (K-like) and hidden-strange (ϕ -like) hadrons. The main problem here is the effect of the substantial 'current mass' of the s-quark (~ 150 MeV) which prevents a literal application of the above DB χ S mechanism as employed for ud-quarks. To adapt the DB χ S mechanism to K- and ϕ -mesons, we use a simple expedient which effectively "reduces" their problem to that of non-strange mesons, viz., we assume that in the first (additive) approximation the K- or K^* - masses exceed those of the corresponding non-strange mesons by the difference between current masses of s- and ud- quarks. This difference in turn can be estimated, a la Politzer [6], as the difference (150 MeV) between the corresponding constituent values, viz., $m_s = 415$ MeV, $m_{ud} = 265$ MeV which fit the spectra [8]. Therefore it is first necessary to make an adhoc subtraction of m_c (≈ 150 MeV) from their actual masses before these resonances can be subjected to the above DB χ S treatment (with $m_q = 265$ MeV and $m_f = 176$ MeV). Similarly for ϕ -like mesons, a prior subtraction of $2 m_c$ (≈ 300 MeV) is necessary for all of them.

The results for K- and ϕ -like mesons are shown in table 2, where a sliding value of m_c (150-200 MeV) has been used to keep a continuous tab on the variation of the total widths with this parameter which is (expectedly) smooth. The calculation of table 2 is otherwise on identical lines to those of table 1. There are now more possibilities of mixing among kaon states than among the corresponding non-strange mesons, since S and C are no longer good quantum numbers, though in table 2 we have not included such effects in view of our semi quantitative description of these cases, which does not make such exercise worthwhile. Apart from this one expects 4 quark components to compete more often with the $q\bar{q}$ component in

several of these cases, especially those of $J = L = 1$, viz., $K^*(1410)$, $K_0^*(1430)$, $K^*(1680)$, $f_0(1400)$, for which the results seem to be particularly bad. Nevertheless the general pattern of the results does seem to indicate a fair degree of overlap with the data upwards of $N=1$ states, just as for table 1, despite the fact that the theoretical status of table 2 is weaker than that of table 1.

In achieving this limited agreement, the interference between A_{NP} and A_P has proved crucial. In this respect, the smaller contribution of A_{NP} than that of A_P may be viewed in the perspective of momentum variation of the wave function in the respective momentum domains, despite a naive expectation to the contrary born out of the relative roles of the 'soft' (h.o.) and 'hard' (o.g.e.) domains in respect of the mass (M) calculations [8,13]. Indeed, these two domains seem to play complementary roles for the width calculations where the large momentum dependence of the wave function plays a crucial role, vis-a-vis the mass calculations [8,13], where this effect is much less important. Nevertheless the relative signs of A_{NP} and A_P are correctly given by Eqs. (23-25) and correspond to those of the mass (M) calculations [8,13] as well, and we would like to emphasize the crucial role of their (destructive) interference in producing the agreement shown in table 1. To check this point, the separate columns Γ_{NP} and Γ_P have also been shown to bring out their individual contributions which are seen to be "too low" and "too high" respectively for the data.

To summarise, we have tried to give a rather simple minded description of the total widths of the $q\bar{q}$ hadrons through a coherent interference of the two diagrams corresponding to soft and hard gluon-exchange mechanisms respectively, using an inclusive mechanism borrowed from DIS and heavy to light transitions through the ansatz of a small time zone during which the quark pair from the hadron is in a quasi-free state before eventual hadronization. This mechanism, symbolized by fig.3(a,b) for the OZI-allowed light $\rightarrow$ light transitions, does not have a counterpart in the corresponding gluon emission mechanisms, fig. 1(b,c), for OZI forbidden heavy $\rightarrow$ light transitions. In this respect a non-perturbative mass formula of the

type (17), incorporating the dynamical breaking of chiral symmetry (inherent in a 4-fermion interaction Lagrangian mediated by a vector-like exchange) through an interplay of SDE and BSE has been the key ingredient[10] . We are not yet in a position to offer a more convincing support for the above ansatz , and so have had to rely on its observational effects . To that end , we have insisted on the precalibration of the mass function from spectroscopy [8,13], and also checked that it yields a parameterfree value $(266 \text{ MeV})^3$ for $\langle u\bar{u} \rangle_0$ in accord with QCD sum rules [16]. We have found that, except for some $L \leq 1$ cases which are susceptible to hadronic selection rules in varying degrees, this simple model provides a rather good description of the quantitative pattern of the observed widths all the way upto $J = 5$, an agreement in which the (destructive) interference of the non-perturbative amplitude A_{NF} and the perturbative amplitude A_P has proved quite important. Applications to baryonic decays are under way .

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References :

1. Particle Data Group, Phys.Rev.D45-Part II, June 1992
2. R.P.Feynman, M.Kislinger and R.Ravndal, Phys.Rev.D3, 2706 (1971)
3. Y.Nambu and G.Jona Lasino, Phys.Rev.122, 345 (1961)
4. S.Adler and A.C.Davies, Nucl.Phys.B244, 469 (1984)
5. R.Delbourgo and M.D.Scadron, J.Phys.G5, 1621 (1979)
6. H.D.Politzer, Nucl.Phys.B117, 397 (1976)
7. K.Lane, Phys.Rev.D10, 2605 (1974)
8. K.K.Gupta et al, Phys.Rev.D42, 1604 (1990);
D.S.Kulshreshtha and A.N.Mitra, Phys.Rev.D37, 1268 (1988)
9. A.N.Mitra and S.Bhatnagar, Intl.J.Mod.Phys.A7, 121 (1992)
10. A.N.Mitra and B.Mitra, Intl.J.Mod.Phys.A9, 915 (1994).
11. G't Hooft, Nucl.Phys.B75, 461 (1974) ;
E.Witten, Nucl.Phys.B160, 57 (1979).
12. A.N.Mitra and D.S.Kulshreshtha, Phys.Rev. D26, 3231 (1982) ;
D28, 588 (1983).
13. A.Mittal and A.N.Mitra, Phys.Rev.Lett. 57, 290 (1986).
14. E.g. M.Shifman, Nucl.Phys. B173, 13 (1980).
15. E.g. A.V.Radyushkin, Phys. Lett. B271 , 218 (1991).
16. V.L.Chernyak and A.R.Zhitnitsky, Phys.Rep.112, 174 (1984)
17. S. Bhatnagar et. al. Phys. Lett. B263, 485 (1991).
18. V.Elias and M.D.Scadron, J Phys.G14, 1175 (1988)
19. A.N.Mitra Zeits.f.Phys. C8 , 25 (1981)
20. J.Mandula et al , Ann.Rev.Nucl.Sci. 20 , 289 (1970)

Fig. 1a : Decay of a hadron via DIS process.

Fig. 1(b,c) : Diagrams for heavy $\rightarrow$ light flavour transitions of $C = \pm 1$ quarkonia via gluon emissions

Figs. 2(a,b,c...) : Decay of a light quarkonium ($q\bar{q}$) into similar hadronic species. The quasi-free state of the quark pair is represented between the vertical dotted lines. For details see text.

Figs. 3(a,b) : Light $\rightarrow$ Light flavour transitions via soft (dotted line) and hard (coiled line) gluon exchange mechanisms respectively.

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Fig. 1(a,b,c)

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Fig. 2 (a, b, c)

Table 1

Total widths (in Mev) of non-strange hadrons with specified [10] values of β^2 (in GeV^2). The widths Γ_{NP} , Γ_P , Γ_t are the contributions of Fig. (3a), (3b) and their resultant respectively.

Meson	$I^G J^{PC}$	β^2	N_H^{-2}	Γ_{NP}	Γ_P	Γ_t	$\Gamma(\text{expt.})$
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V_J mesons :

$\rho(775)$	$1^+ 1^{--}$	0.0697	.0511	83	336	87	153 ± 3
$f_2(1270)$	$0^+ 2^{++}$	0.0935	.1579	26	377	205	185 ± 20
$a_2(1320)$	$1^- 2^{++}$	0.0959	.1689	23	416	244	110 ± 5
$\omega_3(1670)$	$0^- 3^{--}$	0.1130	.3642	12	336	221	166 ± 15
$\rho_3(1690)$	$1^+ 3^{--}$	0.1140	.3727	12	344	230	215 ± 20
$f_4(2050)$	$0^+ 4^{++}$	0.1313	.7270	6	315	232	204 ± 13
$\rho_5(2350)$	$1^+ 5^{--}$	0.1456	1.2319	4	269	206	
$f_6(2510)$	$0^+ 6^{++}$	0.	1.7479	4	213	158	

B^J mesons :

$h_1(1170)$	$0^- 1^{+-}$	0.0886	.2001	36	346	159	335 ± 26
$b_1(1235)$	$1^+ 1^{+-}$	0.0918	.2200	31	384	198	150 ± 10
$\pi_2(1670)$	$1^- 2^{+-}$	0.1130	.5780	13	357	235	250 ± 20

S_0 mesons :

$f_0(975)$	$0^+ 0^{++}$	0.0791	.1053	29	133	38	3^A
$a_0(980)$	$1^- 0^{++}$	0.0793	.1059	29	134	39	57 ± 11

A_1 meson :

$a_1(1260)$	$1^- 1^{++}$	0.0930	.1257	21	287	154	$300-600$
$f_1(1285)$	$1^+ 1^{++}$	0.0942	.1306	19	295	164	24 ± 3

Table 2

Total widths (in MeV) of strange and hidden strange hadrons. The entries Γ_{NP} , Γ_P , Γ_t are the contributions of Fig. (3a), (3b) and their resultant respectively, as in table 1. The subtractions $m_c = 150-200$ MeV from the kaon masses, and $2 m_c = 300-400$ MeV for the ϕ -like masses, have been made beforehand to bring their width calculations in line with those of non-strange mesons, table 1 (see text for explanation). The indicated range of values for the different entries, corresponds to the range 150-200 MeV for m_c . (The quantum numbers G and C are meant only for the ϕ -like hadrons.)

Meson	$I^G J^{PC}$	Γ_{NP}	Γ_P	Γ_t	$\Gamma(\text{expt.})$
V_J mesons :					
$K^*(892)$	$1/2 \ 1^-$	90-105	267-236	49-30	50
$\phi(1020)$	$0^- \ 1^{--}$	95-114	252-207	40-16	4.43
$K_2^*(1430)$	$1/2 \ 2^+$	25-29	383-338	211-169	109 ± 5
$f_2'(1525)$	$0^+ \ 2^{++}$	29-37	334-268	166-106	76
$\phi_3(1850)$	$0^- \ 3^{--}$	16-21	277-243	159-122	87 ± 6
$K_3^*(1780)$	$1/2 \ 3^-$	13.5-15	303-287	188-170	164 ± 17
$K_4^*(2045)$	$1/2 \ 4^+$	9-10.5	244-224	158-137	198 ± 30
$f_4(2300)$	$0^+ \ 4^{++}$	7-9	284-246	201-160	150-200
B_J^J mesons :					
$K_1(1270)$	$1/2 \ 1^+$	40-45	289-264	113-91	90 ± 20
$K_2(1270)$	$1/2 \ 2^-$	15-17	320-302	197-177	136 ± 18
S_J mesons :					
$f_0(1400)$	$0^+ \ 0^{++}$	19-24	127-101	47-29	150-400
$K_0^*(1430)$	$1/2 \ 0^+$	10-12	152-143	84-72	287 ± 23
$K^*(1680)$	$1/2 \ 1^-$	67-73	105-85	59-41	323 ± 110
A_1 meson :					
$K_1(1400)$	$1/2 \ 1^+$	21-25	278-256	145-121	174 ± 13
$f_1(1510)$	$0^+ \ 1^{++}$	24-33	260-228	125-87	35 ± 15