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A Minimal Dynamical Scheme for "Understanding" Hadronic Widths

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ABSTRACT

A quark model scheme of 'least Dynamics' is proposed for an inclusive determination of the total width of hadrons (meson and baryon) via the following mechanism : There exists a short (breathing) mode for the temporary dissociation of a hadron into a pair of "quasi-free" quarks and/or diquarks with significantly lower masses than their normal (constituent) values . The further ansatz of eventual hadronization (with certainty) yields the requisite widths from the (first stage) dissociation probability itself. The coupling scheme proposed for the meson - quark vertex function is based on the general principles of (broken) chiral symmetry as well as the conservation of isospin, baryonic charge and hypercharge with respective 'gauge carriers' ρ , ω , ϕ . A similar construction holds for baryon couplings to quark (q)- diquark (D) pairs , with the same overall coupling constant (g_s) for both systems. Its tentative identification with the QCD value yields a surprisingly good pattern of agreement with data for both hadron types, using just two inputs (m_{ud} and m_s), with the corresponding diquark masses determined additively.

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1. Introduction and Motivation :

The total width (Γ) of a hadron is an intrinsic property on par with its mass (M), as can be seen from the listings in the PDG [1] tables, yet it has long been customary [2] to attempt its saturation via different 'exclusive' branching modes only. On the other hand an increasing tendency in the PDG tables over the years has been to list with some definiteness only the (M, Γ) values, while the different branching modes (especially for the meson listings) are generally too qualitative and too inadequate to allow any meaningful saturation via exclusive modes only. This already makes an apriori case for a direct inclusive determination of Γ , were it not for the (formidable) problem of quark confinement which rules out a literal implementation of such a program without proper **theoretical** safeguards. Nevertheless the challenge is real, and calls for alternative strategies. In this context it is worth noting a recent interesting finding [3] of a highly suggestive **parametrization** of the form $\Gamma / M = n / 2^n$ between the masses and total widths of hadrons not only in the strong interaction sector (mostly resonances) but over the entire PDG listings (and more). Though efforts of this nature date back to Nambu [4], the wide span of these findings [3] calls for a closer scrutiny, especially in the strong interaction sector which represents the central motivation for this investigation of total widths of hadrons, though strictly within the orthodox theoretical framework.

1.1 A simple mechanism for total widths :

The problem of total widths of light qqbar meson resonances was recently addressed [5] in the context of a Bethe - Salpeter (BSE) cum Schwinger - Dyson (SDE) model of an effective qqbar Lagrangian (gluon-exchange-like) with current quarks which guarantees chiral symmetry at the input Lagrangian level, but breaks the symmetry dynamically (D χ SB effect) through the solution of the SDE, generating in the process a momentum -dependence for the quark's mass function $m(p)$. This well known result [6] was merely implemented in the context of a specific BSE model [7] attuned to spectroscopy. The latter also yielded an explicit non-perturbative formula for $m(p)$ [8] through the use of the result [6] that in the (chiral) limit of vanishing pion mass the pion-qqbar vertex function reduces to $m(p)$, so that $m(p)$ could be inferred from the pion wave function derived from BS-dynamics as formulated in [7].

Using this machinery the problem of total widths [5] could be addressed as follows. One first notes that unlike the case of DIS (where a hard e.m. 'kick' is available for dissociation into quarks), or the case of heavy (QQbar) transitions to lighter (qqbar) species which occur via gluonic links, no such simple mechanism exists for light qqbar transitions to the same species. This necessitated an additional assumption for the transition process, namely, there exists a short hiatus period (a sort of "breathing mode") during which a meson dissociates into a pair of quasi-free quarks which must eventually 'hadronize with

certainty'. The transition to this state may be brought about via the exchange of a gluon (hard and soft) , and the resulting masses of the quasi-free quarks are determined by the appropriate momentum of the mass function $m(p)$ to give a much smaller value (176 MeV) [5] than the 'constituent' value (265 MeV) [7] which corresponds to $p = 0$. With such a 'quasi-free' quark mechanism, the calculation of total widths reduces to that of the first stage probability only, requiring merely the hadron- $q\bar{q}$ vertex function with gluon exchange (calculable in the BSE framework). The results which were otherwise parameterfree showed a reasonable but limited pattern of agreement with data for excitations upwards of $L = 1$, while those with $L \leq 1$ were susceptible to hadronic selection rules which this formalism could only recognize but was not equipped for, since the flavour dependence of the $q\bar{q}$ Lagrangian had not been incorporated in the formalism . On the other hand the results did seem to bear out the basic viability of the suggested mechanism of (temporary) dissociation into a pair of quasi-free quarks before eventual hadronization. (Other details like vanishing widths for $N_c = \text{infinity}$ were trivially satisfied).

1.2 Quark-diquark description :

In this paper we wish to pursue the consequences of this basic mechanism (of temporary quasi-free dissociation before eventual hadronization with certainty) on both types of hadrons in some detail so as to permit a more serious (quantitative) comparison with data [1]; this necessitates the incorporation of the major

hadronic selection rules within the formalism itself. Partly for this reason, but mainly to facilitate a more direct comparison of the basic assumption of 'quasi-free dissociation' with the data without the intervention of too specific a model [5-8], it is useful (and instructive) to consider a simpler approach which permits the incorporation of the essential features of the dynamics, including the role of the major symmetry - breaking effects, without having to go into any 'excessive details thereof'. (For this reason we have preferred using the word dynamical scheme, instead of a 'dynamical model', in the title of this paper). In the spirit of such simplification it should be legitimate to employ the concept of the diquark (D) [9] to compactly characterize the hadronic structures under study. In this picture a meson is a linear combination of $q\bar{q}$ and DD^* states, while a baryon is dominantly a qD state with possible admixtures of higher Fock components. In this way one formally avoids a 'three-body' description. [Though a DD^* component in a meson amounts to a $q^2\bar{q}^2$ configuration, we refrain from using a $(q\bar{q})^2$ [two-meson (color-singlet) description] for fear of double counting of the available d.o.f.'s]. As to the masses of these objects, since we shall need them in their 'quasi-free' states, it is useful to continue to employ the values obtained from the more formal dynamical model attuned to spectroscopy [5,7], namely, $m(q) = 176 \text{ MeV}$ for non-strange, and $m(q) = 176 + 150 = 326 \text{ MeV}$ for strange quarks, the figure of 150 being the 'normal' difference of their 'current' masses (which also checks with the

spectroscopic determinations [7] of their 'constituent' values, viz., 265 and 415 respectively, under Politzer additivity [10]). As to the corresponding diquark masses we again consider the simplest possibility, viz., $m(D) = 2 \times 176 = \underline{352 \text{ MeV}}$ (non-strange); $m(D) = 2 \times 176 + 150 = \underline{502 \text{ MeV}}$ (strange), etc., all under simple additivity, so as to avoid fresh parameters without good reason.

1.3 Gauge couplings and Partial Symmetry :

For the essential dynamics, one may profitably take over some valuable results from yesteryear and adapt them to the couplings of hadrons to quarks and diquarks. One outstanding result due to Sakurai [11] concerns the gauge couplings of ρ , ω and ϕ mesons to ensure the conservations of isospin, baryonic charge and hypercharge respectively. Another highly compact expression of a comprehensive set of couplings under broken chiral symmetry at the hadronic level comes from the phrase "Partial Symmetry" due to Schwinger [12] who was able to arrange the (scattered) information on various current-algebraic results under the umbrella of a master 'meson-matrix' [12]. Through a simultaneous use of these two principles it becomes possible to organize the entire coupling structure of mesons to $q\bar{q}$ and DD^* systems under a single overall coupling constant (g_s). And a trivial extension of the same ideas also provides an equally complete basis for the corresponding couplings of baryons to qD systems, with the same overall coupling constant (g_s) for both coupling types.

Now since most of the hadrons listed in the PDG tables [1] are mere L- and/or radial- excitations of the corresponding ground states , the same principle of Partial Symmetry [12] can be quite straightforwardly extended [13] to the excited states too , with little change in the basic structure (see Sec.2 below). It is now tempting to identify, at least tentatively, the quantity g_s with its standard QCD value, viz., $g_s^2/4\pi = \alpha_s$. Though there is as yet no formal basis for such optimism , we shall nevertheless make such tentative identification which is in any case in conformity with the spirit of a dynamical 'scheme' (as distinct from a regular dynamical 'theory'), to see if a minimal degree of order exists in the data [1] on rather widely varying hadronic widths. Finally, for the "form - factors" for the hadron-quark (diquark) vertices , which could be quite important for the excited hadrons, we again take the "path of least dynamics" which essentially amounts to replacing the various 'momentum 4-vectors' appearing in the (multiple index) tensorial couplings [13], by the corresponding 'velocity 4-vectors' [14] (whose use is now quite fashionable in the context of 'heavy flavour symmetries' [15]).

1.4 Contents of the Paper :

With the broad contours of our dynamical scheme so defined, the next few sections are devoted to a summary of its quantitative formulations so as to facilitate the explicit calculations as well as the numerical results. Sec.2 summarises (in 3 parts)

the structure of the "Partial Symmetric" meson matrix [12] and its L-excited generalization [13] at the $q\bar{q}$ and DD^* levels of couplings. In view of the existence of a large variety of DD^* pairs, some further criteria (in the form of suitable selection rules for good G-parity mesons only) to delimit their choice are employed in terms of some standard quark-model rules, with diquark [9] units, as guiding principles. Some essential steps for the actual construction of these couplings are briefly recalled [13] in Appendix A. Sec.3 sketches the analogous derivation of the coupling structures for the dissociation of a baryon into a quasi-free qD pair, with its qD content known from the $SU(6) \otimes O(3)$ quark model providing the necessary rationale for the baryon- qD couplings. [The results on decay widths for both types of hadronic couplings are collected in Appendix B]. Sec.4 presents the results, separately for mesons and baryons of different types (non-strange, strange, etc.). It is found that the general pattern of the width variations is indeed well reproduced over the entire PDG tables for both types of hadrons. The main lessons from these results which are expressed in the form of simple analytical formulae for the ratio $R \equiv \Gamma/M$, are examined vis-a-vis the central features of the theoretical postulates. Sec.5 summarises our conclusions.

2. Meson Matrix for qqbar and DDStar Couplings :

The content of Partial Symmetry [12] may be expressed as a 6x6 hermitian matrix (for three flavours) made out of the positive parity field components of Vector (rho, omega and phi) and Pseudoscalar (pion, kaon and eta) mesons whose couplings to a $q\bar{q}$ system survive in the limit of small momentum transfer. Taking account of the mass differences among vector mesons (which determine the ratio of their magnetic to charge couplings), Schwinger gave an explicit formula for their e.m. couplings to baryons [12]. The same structure may be adapted to the $q\bar{q}$ level of strong interactions, separately for 'isospin', 'baryonic charge' and hypercharge, following the 'gauge prescription' of Sakurai [11] but otherwise within the Schwinger framework [12] itself.

2.1 qqbar Couplings of L=0 Mesons :

The ground state meson matrix for $q\bar{q}$ couplings in the limit of small 3-momentum transfer k_i reads in momentum space as [13,16] :

$$MM(q\bar{q}) = (g_s/2) \sum_{a=0}^8 \lambda_a [i(\sigma_i k_i \pi^a - \epsilon_{ijm} \sigma_i V_j^a k_m) / M_V^a + V_0^a + \sigma_i A_i^a] \quad (1)$$

Here 'a' is the U(3) flavour index ; π^a , V_μ^a , A_μ^a are the generic names for pseudoscalar, vector (mass M_V^a) and axial vector fields.

In the traditional picture (with L=0 mesons only), the λ -matrices given above for ω and ϕ are to be read in the sense of 'ideal mixtures' while the others retain their usual meaning. Now to

couple these fields to the $q\bar{q}$ system their "strengths" have to be so arranged as to give expression to our stated program of gauge couplings of ρ to isospin, ω (and f_2) to baryonic charge, and ϕ to hypercharge [11]. For the $L=0$ mesons, however, we wish to play safe and merely assume that their 'strong' couplings follow exactly the pattern of their electromagnetic charge couplings, which just amounts to incorporating Schwinger's e.m. substitutions [12] in the definitions of the λ -matrices. In this respect the ρ coupling is already normalized as it stands, i.e., $\lambda_a = \tau_a$ for $a = 1, 2, 3$. the same goes for K and K^* λ 's ($a=4-7$). However for $a=0$ and 8 , the e.m. substitution rule amounts to using the following definitions of the λ -matrices associated with ω and ϕ :

$$\lambda_\omega = \frac{1}{3} \begin{vmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 0 \end{vmatrix} ; \quad \lambda_\phi = -\frac{2}{3} \begin{vmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{vmatrix} \quad (2)$$

while all other λ -matrices in (1), including the unnatural parity (π -like) mesons, retain their usual meaning, thus automatically defining their strong couplings relative to the V -mesons. [The strength of the ω coupling to $q\bar{q}$, given by eq.(2), may be compared to its isoscalar charge coupling to e^+e^- ; and that of ϕ to its hypercharge coupling to e^+e^- With the correspondence " e " $\leftrightarrow$ " g_s "].

With these specifications, the Lorentz-boosted $q\bar{q}$ couplings of the $L=0$ hadrons have the resultant form :

$$MM(q\bar{q}) = (g_s/2) \sum_a [\pi^a \lambda_a i\gamma_5 \gamma \cdot P/M_V + i\gamma_\mu V_\mu^a \lambda_a + i\gamma_5 \gamma_\mu A_\mu^a \lambda_a] \quad (3)$$

sandwiched between the $q\bar{q}$ 4-spinors; here the axial partners A_μ^a of V_μ^a have been shown to appear with equal strengths appropriate for chiral symmetry. (The other alternative to generate the axial couplings is via the $L=1$ excitation mode of an $S=1$ $q\bar{q}$ system which gives $J = 2, 1, 0$ respectively; we adopt the former for $I=1$ and the latter for $I=0$). Note that in (3) we have left out the Lorentz-boost of the magnetic term. This is just because the coupling is at the quark level (no anomalous magnetic moment!), the γ_μ term itself incorporates the effect of its (normal) m.m., so that any extra contribution would amount to overcounting. This does not however preclude the use of the magnetic term for the coupling of all hadrons (including $L=0$) to quark composites like DD^* (see Appendix A for details).

2.2. $q\bar{q}$ Couplings of L-excited Mesons :

Before using eq(1) for the general case, it is necessary to incorporate the 'quark recoil effect' [17] on the one hand, and the 'convective current' [2] on the other. Both these effects are particularly useful in giving a self-consistent shape to the Lorentz - boosting of these couplings [13] before they can be employed for the actual calculations. This amounts to the replacements [13]

$$k_i \rightarrow P_i - M p_i^{(1,2)} / \omega_{(1,2)} ; V_0 \rightarrow V_0 - 2 p_j^{(1,2)} V_j / M \quad (4)$$

where (M, P_i) are the mass and 3-momentum of the meson, taken as 'radiation quantum' [18] (at rest), and $(\omega^{(1,2)}, p_i^{(1,2)})$ are the energy-momentum of the quark /antiquark interacting with the former in turn. These are particularly necessary for the couplings of hadrons to the DD^* pairs (see sec.2.3 and Appendix A).

For the $q\bar{q}$ couplings of L-excited mesons, we first need to modify the definition (2) of the λ_ω and λ_ϕ matrices so as to conform directly to the roles of the associated mesons ω^L and ϕ^L as gauge carriers of baryonic charge and hypercharge respectively, with full strength g_s instead of $g_s/2$ as above for $L=0$. (The fact that we could not implement this requirement all the way down to the $L=0$ level reflects a limitation of our 'scheme' vis-a-vis the data). To that end, the appropriate definitions of λ_ω and λ_ϕ are

$$\lambda_\omega = \frac{2}{3} \begin{vmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{vmatrix} \quad ; \quad \lambda_\phi = \frac{2}{3} \begin{vmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -2 \end{vmatrix} \quad (5)$$

However we shall continue to regard the resonances ω_J and ϕ_J as ideally mixed for purposes of our calculations .

For the L-excited states we must use the multiple tensor formalism [19] with multiple derivative couplings [20]. To that end we define the relative 4-momentum Q_μ of the two (outgoing) quarks (m_1, p_1) and (m_2, p_2) , which for equal masses is $(p_{1\mu} - p_{2\mu})/2$ (for unequal masses, see eq.(A.2) of Appendix A). Using this quantity it is convenient to define the Q-contracted tensors [5]

$$B^L(Q) = Q_{\mu_1} Q_{\mu_2} \dots Q_{\mu_L} B^L_{\mu_1 \mu_2 \dots \mu_L}(P) ; Q.P = 0. \quad (6)$$

$$V^J_\mu(Q) = Q_{\mu_1} Q_{\mu_2} \dots Q_{\mu_L} V^J_{\mu \mu_1 \mu_2 \dots \mu_L}(P) ; J = L + 1. \quad (7)$$

$$S^J(Q) = Q_{\mu_1} Q_{\mu_2} \dots Q_{\mu_J} S^J_{\mu_1 \dots \mu_J}(P) \equiv Q_{\mu_1} S^J_{\mu_1}(Q) ; J = L - 1 \quad (8)$$

$$A^J_\mu(Q) = \gamma_5 V^J_\mu(Q) ; \quad (9)$$

As the notation indicates, these symbols stand for the $q\bar{q}$ couplings of L-excited pseudoscalar (b_1 -like), vector (f_2 -like), scalar (a_0, f_0 -like) and axial-vector (a_1 -like) mesons, with the further proviso (as noted above and specified by eq(9)) that a_1 is paired with rho, etc. The next point concerns the relative strengths of these $q\bar{q}$ couplings for a given L-value. For the natural parity ($S = 1$) mesons this is given by the coefficients of the C.G. series for $T(L) \otimes S(1)$ which would normally be dependent on the m (magnetic) quantum numbers and hence inconvenient for relativistic generalizations. A much better alternative is to employ a direct tensorial reduction [13]

$$V_i \otimes B^L(Q) = V_i^{L+1}(Q) + i \eta(L) \epsilon_{ijm} Q_m A^L_j(Q) + \\ + \zeta(L) (2 Q_i S^{L-1}(Q) + (L-1) Q^2 S^{L-1}_i(Q)) ;$$

$$\eta^2(L) = 2L^2/(L+1)^2 ; \zeta^2(L) = L^2 / [(L+1)^2 (4L^2-1)] \quad (10)$$

which is a 3D adaptation, in the notation of eqs.(6-9), of a derivation given in appendix D of ref.[13]. (Some marginal errors

in the derivation of these coefficients have been corrected in (10) }. The normalizations of these Q-contracted tensors are as follows [5,13,19] :

$$\langle B^J(Q)^* B^J(Q) \rangle_{AV} = Q^{2J} / p(J); \quad p(J) = (2J+1)!! / (J)! \equiv \xi^2(J) \quad (11)$$

$$\langle V_\mu^J(Q)^* V_\nu^J(Q) \rangle_{AV} = Q^{2J-2} \{ (J+1) \Theta_{\mu\nu}(P) + (J-1) \hat{Q}_\mu \hat{Q}_\nu \} / 2p(J) \quad (12)$$

$$\Theta_{\mu\nu}(P) = \delta_{\mu\nu} + P_\mu P_\nu / M^2; \quad \hat{Q}_\mu \hat{Q}_\nu = Q_\mu Q_\nu / Q^2 \quad (13)$$

The resultant $q\bar{q}$ coupling of L-excited mesons is now given by

$$\begin{aligned} M_L(q\bar{q}) = & \xi(L) (g_s/2) \sqrt{L} \sum_a [B^{La}(\hat{Q}) \lambda_a i \gamma_5 \gamma_\mu P/M_L + i \gamma_\mu V_\mu^{(L+1)a}(\hat{Q}) \lambda_a + \\ & + \eta(L) i \gamma_5 \gamma_\mu A_\mu^{La}(\hat{Q}) \lambda_a + \zeta(L) (2 i \gamma_\mu Q S^{L-1}(Q) + \\ & + (L-1) Q^2 i \gamma_\mu S^{L-1}(Q)) \lambda_a] \end{aligned} \quad (14)$$

again sandwiched between $q\bar{q}$ 4-spinors ; the λ matrices for $a = 0$ and 8 natural parity states have the same significance as in eq. (5) and the CG coefficients are defined as in (10). The first term which corresponds to the unnatural parity (pi-like) couplings has a dividing factor M_L which corresponds to the mass of the 'associated' $V^{(L+1)}$ member of the nonet. The "velocity" factor [14,15] sitting in front is given by

$$\sqrt{L} = (2 Q/M)^L; \quad Q^2 = \lambda(M^2, m_1^2, m_2^2) / 4 M^2 \quad (15)$$

$\lambda(\dots)$ being the standard triangle function of the arguments.

2.3. Meson DDstar Couplings :

We now indicate briefly the corresponding meson couplings to the

admissible DD^* states on analogous lines to ref.[16] for the Partial Symmetric e.m. couplings for baryonium states. Physically the inclusion of this channel amounts to a recognition of the importance of "four-quark" states in the meson dynamics as has been emphasized by some authors [21]. (Indeed we may as well anticipate that significant contributions to the total widths are going to arise from the DD^* channel). There are now two types of (color 3^*) diquarks, axial vector D_μ and scalar D_s , and their corresponding antiparticles D_μ^* and D_s^* (color 3). For the $SU(3)$ indices, however, it is simplest to use the isovector index D_μ^a or D_s^a ($a = 1, 2, 3$) with the matrix T_a as needed, as well as isoscalar and isospinor (indices suppressed) notations with D_μ or D_s when relevant. We recall in this context that the 21 representation of $SU(6)$ consists of D_μ^a and D_s which make up the baryonic 56 states together with an extra quark [9], while their 'complements' D_μ and D_s^a are needed only for the excited (70) baryons. It is helpful to keep this distinction in mind as a relevant analogy in deciding which pairs of DD^* couple dominantly to a meson of a given L-excitation, and hence as a practical aid in formulating some 'selection rules' for such couplings.

In this spirit we shall tentatively adopt the following selection rules for the DD^* coupling of L-excited non-strange, 'natural parity' mesons of "good" G-parity: For even L the coupling is dominantly to D_μ^a and D_s ; for odd L, their main coupling is to D_μ and D_s^a . The only exception to this rule for G-parity

their respective mass differences without change in the coupling forms. For the rest we keep to the PDG notation [1].

Using the above classification, the full Lorentz structures of a sufficiently representative set of states including their 'flavours' are listed below

$$b_L = 2 g_s W_L \xi(L) \sqrt{L} \epsilon_{\mu\nu\alpha\beta} \epsilon_{abc} D_{\mu a}^*(p_2) B_b^L(\hat{Q}) D_{\nu c}(p_1) P_{\alpha} Q_{\beta} \quad (17)$$

$$K_{Lm} = 2 g_s W_L \xi(L) \sqrt{L} D_s^* \{ D_{\mu} + \tau_a D_{\mu}^a \} K^L + h.c. \quad (18)$$

$$\rho_J = 2 g_s W_L \xi(L) \sqrt{L} [D_{\mu}^*(p_2) \times T_a D_{\nu}(p_1)] \cdot [\rho_{\mu}^{Ja}(\hat{Q}) \times Q_{\nu}] \quad (19)$$

$$\rho_{Jm} = 2 g_s W_L \xi(L) \sqrt{L} \epsilon_{\mu\nu\lambda\sigma} D_s^*(p_2) \rho_{\mu}^{Ja} Q_{\nu} D_{\lambda}^a(p_1) P_{\sigma} + h.c. \quad (20)$$

$$K_J^* = 2 g_s W_L \xi(L) \sqrt{L} (D_{\mu}^*(p_2) \times \{ D_{\nu} + \tau_a D_{\nu}^a \}) \cdot (K_{\mu}^{*J}(\hat{Q}) \times Q_{\nu}) \quad (21)$$

where the isospin dependence of the diquark fields is expressed in a matrix notation with the isovector T_a sandwiched between them, and the $\times$ symbol denotes 4D cross product.

$$\omega_{Js} = 2 g_s W_L \xi(L) \sqrt{L} [D_s^*(p_2) Q_{\mu} \omega_{\mu}^J(\hat{Q}) D_s(p_1)] \quad (22)$$

$$\begin{aligned} \omega_J' = & g_s W_L \xi(L) \xi(L) \sqrt{L} \{ 4 \omega^J(\hat{Q}) D_{\mu}^*(p_2) D_{\nu}(p_1) [Q \cdot \hat{Q} \delta_{\mu\nu} - Q_{\mu} \hat{Q}_{\nu}] \\ & + 2 J (Q_{\nu} \times \omega_{\mu}^J) \cdot (D_{\mu}^* \times D_{\nu}) \} \quad ; \end{aligned} \quad (23)$$

The last two couplings refer to $J=L+1$ and $J=L-1$ respectively.

since in this case there is essentially one Lorentz-structure. A second set of couplings of these particles to $D_\mu D_S^*$ states will be denoted with an extra suffix 'm' in (16), e.g. b_{Lm} , k_{Lm} , etc.

For the natural parity states, the flavour structures can again be given by a similar notation, but now a further classification is necessary to distinguish among the different Lorentz- types. Here again we go mostly by the PDG classifications. Thus ρ_J will stand for the $D_\mu D_\nu^*$ - coupling for the odd-J rho-series ($J=L+1$), while ρ_{Jm} and ρ_{JS} will denote their $D_\mu D_S^*$ and $D_S D_S^*$ couplings respectively. Similarly a_J, a_{Jm}, a_{JS} will denote the corresponding couplings of the even J series of $I=1$. (For $J=L-1$, $I=1$ series the above notation suffices.) For the DD^* couplings of isovector axial (A^J) mesons ($J=L$) the corresponding notations are $\{\tilde{a}_J, \tilde{a}_{Jm}\}$, since the $D_S D_S^*$ coupling is zero.

For the isoscalar mesons of $J = L+1$, the even-L and odd-L DD^* couplings are distinguished by the symbols ω_J/ϕ_J and f_J/f_J' respectively, while the Lorentz structures for $D_\mu D_S^*$ and $D_S D_S^*$ couplings are distinguished by the extra suffixes m, s as above. For the K^* -series with $J=L+1$ the obvious notation is $\{k_J^*, k_{Jm}^*, k_{JS}^*\}$, the extra star being used to emphasize its natural parity status. For the corresponding $J=L$ states (axial) the analogous notation is $\{\tilde{k}_J, \tilde{k}_{Jm}\}$. Finally for radial excitations of the hadrons we shall take only their basic (J,L) values, so that their radial quantum numbers will be reflected only in

3.1 qD Content of Baryons :

As to the essential notations for the various spinors, the FDG classification [1] for the various states already provides the flavour basis with N , Δ^a , Λ , Σ^a standing for isospinor, iso-Rarita-Schwinger, isoscalar and isovector respectively ; 'a' being the isovector index as in the meson case. To incorporate the spin d.o.f., we first note [9] that 56 states are made up of D_μ^a and D_s diquarks, while the 70 states also include the complementary combination (D_μ, D_s^a) , in the notation of Sec.2. To see the connection between the qD and qqq descriptions more precisely, it is useful to recall the qqq structure [2,17], in the notation [13,17] of Verde [22], of the non-strange baryons :

$$\begin{aligned} |N\rangle_{56}^d &= [\chi' \phi' + \chi'' \phi''] \psi^s / \sqrt{2} ; \\ |\Delta\rangle_{56}^q &= \chi^s \phi^s \psi^s ; \quad |\Delta\rangle_{70}^d = [\chi' \psi' + \chi'' \psi''] \phi^s / \sqrt{2} \\ |N\rangle_{70}^d &= [(\chi' \phi'' + \chi'' \phi') \psi' + (\chi' \phi' - \chi'' \phi'') \psi''] / 2 ; \\ |N\rangle_{70}^q &= \chi^s [\phi' \psi' + \phi'' \psi''] / \sqrt{2} ; \end{aligned} \quad (26)$$

where χ, ϕ, ψ denote the spin, flavour, orbital d.o.f.'s respectively. The symbols 'd' and 'q' denote spins 1/2 and 3/2 respectively. To identify the qD contents of these wave functions in a Lorentz invariant form, it is useful to note the following correspondence:

$$\chi' \phi' \leftrightarrow D_s ; \chi'' \phi'' \leftrightarrow i \gamma_5 \gamma_\mu D_\mu^a \tau_a ; \chi^s \phi^s \leftrightarrow D_\mu^a \epsilon_\mu^a ; \quad (27)$$

$$\chi' \phi'' \leftrightarrow D_s^a \tau_a ; \chi'' \phi' \leftrightarrow i \gamma_5 \gamma_\mu D_\mu ; \chi^s \phi' \leftrightarrow D_\mu \epsilon_\mu ; \chi^s \phi'' \leftrightarrow D_\mu^a \tau_a \epsilon_\mu \quad (28)$$

Here for simplicity we have suppressed a basis spinor symbol Ψ on the RHS of (27-28) for all states. However the additional spin and

$$\tilde{a}_J = g_S W_L \xi(L) \eta(L) \sqrt{L} D_\sigma^*(p_2) \epsilon^{\mu\nu\rho\sigma} T^C A_\mu^{JC}(\hat{Q}) P_\nu D_\rho(p_1) \quad (24)$$

$$\tilde{a}_{Jm} = g_S W_L \xi(L) \eta(L) \sqrt{L} D_\sigma^*(p_2) T^C [A_\mu^{JC} D_\mu - \hat{Q} \cdot A^{JC} \hat{Q} \cdot D] + \text{h.c.} \quad (25)$$

The decay widths for representative cases are summarized in Appendix B in conjunction with table 1.

3. Baryon Couplings to Quark-Diquark Pair :

We now turn to the baryon-qD coupling problem for ground and L-excited baryons, ignoring the effects of higher Fock components. The basic couplings will appear in the general form $\bar{q}(p_1) D^*(p_2) \Psi(P)$ after inserting Dirac matrices appropriate for the Yukawa interaction of a scalar or axial vector with a pair of unequal spinors q and Ψ . It is first necessary to get the 'color' factor right, since the diquark is merely a convenient artifact for representing the effect of three quarks. And since the choice of two out of three quarks can be made in three distinct ways, there must be a factor of $\sqrt{3}$ in the resultant probability - hence a factor of $\sqrt{3}$ in the transition amplitude - after which we may forget about the color d.o.f. As to the rest we shall again keep to the minimal dynamical level by choosing the $SU(6) \otimes O(3)$ framework of spin-flavour-orbital d.o.f.'s in a Lorentz-boosted form. In this respect we shall follow closely the procedure outlined in Sec.2 for the meson case.

$$|\Lambda\rangle_{70}^{8d,1d} = \phi'(\chi'\psi'' \pm \chi''\psi')/2 ; |\Lambda\rangle_{70}^{8q} = \phi'\psi'\chi^5/\sqrt{2} \quad (30)$$

$$|\Sigma\rangle_{70}^{10d,8d} = \phi''(\chi'\psi' \pm \chi''\psi'')/2 ; |\Sigma\rangle_{70}^{8q} = \phi''\psi''\chi^5/\sqrt{2} \quad (31)$$

3.2 Baryon qD Couplings :

Using the correspondence (27-28) the qD contents of the various baryon types can be directly adapted to the construction of their couplings in the form $\bar{q} D^* \Psi^J + \text{h.c.}$ E.g., the Δ -D-q coupling can be inferred from its full qD content $D_\mu^a \epsilon_\mu^a \otimes \Psi \psi^5$ as $\bar{q}(p_1) D_\mu^{a*}(p_2) \Delta_\mu^a(P)$, where we have hermitian conjugated the diquark to conserve baryon number, and identified the product $\epsilon_\mu^a \otimes \Psi$ with the indicated Rarita-Schwinger spinor. The orbital function ψ^5 may be taken as absorbed in the overall coupling constant which in the simplest $SU(6) \otimes O(3)$ version of the quark - model may be taken as $g_s/2$, the same as in the meson - $q\bar{q}$ case eq.(1). The quantities q , D and Δ must now be regarded as elementary fields with the indicated 4-momenta and masses m_1, m_2, M . All other $L=0$ decouplet baryons (which have strong decays) may be treated in this way. For example, the Σ^* -D-q coupling has the form $\bar{q}(p_1) D_\mu^{a*}(p_2) \Sigma_\mu^{*a}(P)$, where now Σ^{*a} is an isovector. Taking account of the color factor $\sqrt{3}$, the 56 ($L=0$) decouplet has the coupling :

$$\sqrt{3} (g_s/2) \bar{q}(p_1) D_\mu^{a*}(p_2) \{ \Delta_\mu^a(P) + \Sigma_\mu^{*a}(P) + \tau_a \Xi_\mu^{*}(P) \} + \Omega\text{-term} . \quad (32)$$

For the L -excited states, however, we must take account of the non-trivial overlap effect in the form of the velocity-factor $\sqrt{L}$, again as in the meson - $q\bar{q}$ case. More importantly, the relative coupling strengths of the different (J,L,S) states will be given

by the CG expansion of a direct product of a tensor and a spinor. Using the generic notation $\Psi_{(\mu)}^J$ for multiple - index Rarita-Schwinger spinors, these states will couple to the qD pairs in multiple derivative form [20] which in momentum space is given in terms of the relative 4-momentum Q_μ , Eq.(A.2), for the outgoing qD pair, thus providing the rationale for the velocity form factor v^L . Again it is convenient to use the Q -contracted Rarita-Schwinger spinors $\Psi_\mu^J(\hat{Q})$, or Dirac spinors $\Psi^J(\hat{Q})$, in exact analogy with eq.(10) whose spinor counterpart is now given by the full CG expansion

$$B^L(\hat{Q}) \otimes \epsilon_\mu \otimes \Psi = \Psi_\mu^{L+3/2}(\hat{Q}) + \sum_\mu (J) \quad (33)$$

where the (symbolic) summation $\sum_\mu$ is over $J = L+1/2, L-1/2$ and $L-3/2$. Note that the values $J = L \pm 1/2$ must occur twice each, once as members of quartet (q) spin and once with doublet (d) spin respectively. As in the meson case, the object is to determine these coefficients in a form independent of the magnetic quantum numbers. To that end, we use each term of eq.(10) in conjunction with its spinorial counterpart

$$B^L(\hat{Q}) \otimes \Psi^{1/2} \equiv \Phi^L = \Psi^{L+1/2}(\hat{Q}) + \gamma(L/2+1) i \gamma \cdot \hat{Q} \Psi^{L-1/2}(\hat{Q}) \quad (34)$$

to obtain a total of six terms [13] with transparent meaning which may be inferred by analogy with multipole fields [24]: The four 'electric' terms arising from V^J and S^J with "L" values $L \pm 1$ are identified with the four spin quartet (q) states, while the two 'magnetic' terms arising from A^J with "L" = L are obviously

associated with the two doublet (d) states. Therefore the qD overlaps of these states will now provide a natural classification of the relative strengths of their couplings to the various $56^{q,d}$ and $70^{q,d}$ baryons, to within an overall coupling constant g_5 . (For a fuller derivation of such results, see Appendix D of [13]). The normalizations of the states $\Psi^J(Q)$ and $\Psi_\mu^J(Q)$, with appropriate J-values picked from the the CG series eqs.(33-34), are given , analogously to eqs.(11-13), by [19,13] :

$$\langle \Psi^J(Q) \bar{\Psi}^J(Q) \rangle_{Av} \equiv \frac{1}{2J+1} \Theta(J) \left(\frac{M-i\gamma.P}{2M} \right)$$

$$\Theta(J) = \frac{J+1/2}{2J+2} \hat{\gamma}_\alpha^{(P)} \hat{\gamma}_\beta^{(P)} \frac{1}{p(J-1/2)} \left(\theta_{\alpha\beta}^{(P)} \frac{J+3/2}{2} + \hat{Q}_\alpha \hat{Q}_\beta \frac{J-1/2}{2} \right)$$

(35)

$$\langle \Psi_\mu^J(Q) \bar{\Psi}_\nu^J(Q) \rangle_{Av} \equiv \frac{1}{2J+1} \Theta_\mu^\nu(J) \left(\frac{M-i\gamma.P}{2M} \right)$$

(36)

$\Theta_\mu^\nu(J)$ is a similar expression to $\Theta(J)$ but may be omitted since its indices will get eventually contracted.

We are now in a position to write down the complete Lorentz structures of the qD couplings of the L-excited baryonic towers. To that end, we must construct two distinct covariants Ψ_μ^{Lq} and Ψ_μ^{Ld} corresponding respectively to (q) and (d) baryons , through a simultaneous use of the two CG series (10) and (34) wherein the former is the result of applying (34) on the $V_\mu^{(L+1)}$ and $S^{(L-1)}$ terms of (10), and the latter of doing the same thing on $A_\mu^{(L)}$. The resulting (q) and (d) Rarita-Schwinger covariants are :

$$\begin{aligned}
\Psi_{\mu}^{Lq}(\hat{Q}) = & \Psi_{\mu}^{L+3/2}(\hat{Q}) + \alpha_{1L} [L i \gamma \cdot \hat{Q} \Psi_{\mu}^{L+1/2}(\hat{Q}) + i \hat{\gamma}_{\mu} \Psi_{\mu}^{L+1/2}(\hat{Q})] + \\
& + \alpha_{2L} [2 \hat{Q}_{\mu} \Psi_{\mu}^{L-1/2}(\hat{Q}) + (L-1) \hat{Q} \cdot \hat{Q} \Psi_{\mu}^{L-1/2}(\hat{Q})] \\
& + \alpha_{3L} [2 i \gamma \cdot \hat{Q} \hat{Q}_{\mu} + (L-1) \hat{Q} \cdot \hat{Q} i \hat{\gamma}_{\mu}] \Psi_{\mu}^{L-3/2}(\hat{Q})
\end{aligned} \quad (37)$$

$$\alpha_{1L}^{-2} = [2L+3][L+1] ; \alpha_{2L}^2 = L^2 / [(L+1)^2(4L^2-1)]$$

$$\alpha_{3L}^2 = \alpha_{2L}^2 (L-1)/(2L-1) \quad (38)$$

$$\begin{aligned}
\Psi_{\mu}^{Ld}(\hat{Q}) = & \sqrt{\frac{2}{L+1}} i \hat{Q}_{\nu} \epsilon_{\mu\nu\rho\lambda} \gamma_{\lambda} \gamma_5 [\Psi_{\rho}^{L+1/2}(\hat{Q}) + \{i \hat{\gamma}_{\rho} \Psi_{\rho}^{L-1/2}(\hat{Q}) + \\
& i \gamma \cdot \hat{Q} (L-1) \Psi_{\rho}^{L-1/2}(\hat{Q})\} / \sqrt{L(2L+1)}]
\end{aligned} \quad (39)$$

To complete the coupling structures, the corresponding qD covariants must be attached to (37) and (39) to give finally :

$$\bar{q} D^{*Lq} = \sqrt{3}(g_s/2) \xi(L) \sqrt{L} \bar{q}(p_1) [D_s^{*}(p_2) i \hat{\gamma}_{\mu} + \gamma_5 D_{\mu}^{*}(p_2)] \cdot \Psi_{\mu}^{Lq}(\hat{Q}) \quad (40)$$

$$\bar{q} D^{*Ld} = \sqrt{3}(g_s/2) \xi(L) \sqrt{L} \bar{q}(p_1) [D_s^{*}(p_2) i \hat{\gamma}_{\mu} \gamma_5 + D_{\mu}^{*}(p_2)] \cdot \Psi_{\mu}^{Ld}(\hat{Q}) \quad (41)$$

respectively, where $\hat{\gamma}_{\mu}$ is transverse to the baryon momentum P_{μ} . [We have chosen 'opposite parity' conventions for Ψ^{Lq} and Ψ^{Ld} couplings]. The diquark fields are matrices in isospin space; e.g., for 56 N-states, $D_s \rightarrow D_s I$ and $D_{\mu} \rightarrow D_{\mu}^a \tau_a$, etc. The other symbols $\xi(L)$ and $\sqrt{L}$ are defined in eqs.(11) and (15) respectively. There is now no need to write the coupling structures separately for $56^{q,d}$ or $70^{q,d}$ states since their SU(6) qD contents can be

read from direct inspection of eqs.(26-31).

3.3 Doublet (d) States of $L \equiv 0$:

A closer look at the entire coupling structure for the baryonic towers given by eqs.(37-41) above suggests that the couplings are fully defined for all L -values including the $L=0$ 56^q . The only missing item in the series is the $L=0$ 56^d , since a naive substitution of $L = 0$ in eq.(39) will give a null value. This is perhaps just as well since these **ground state** baryons are stable against strong decay. On the other hand, this $SU(6) \otimes O(3)$ scheme gives zero qD coupling to the radial excitations of 56^d ($L=0$) as well, e.g., the Roper $N(1440) P_{11}$ and similar 70^d states like $N(1710)P_{11}$, which have enough mass for strong decay. It is nevertheless interesting to ask if, without introducing any fresh coupling constant, it is at all possible to address the qD-couplings of such non-orbital excitations of $L^d = 0^d$ states within this very $SU(6) \otimes O(3)$ framework. To that end, consider formally the result of taking the opposite limit, viz., $L = \infty$ in eq.(39), after evaluating their qD dissociation widths. Rather interestingly, this limit gives a finite value, as if to suggest that the route to the $L = 0$ limit lies via $L = \infty$! To exploit this rather intriguing result within our coupling scheme, it is tempting to "propose" that the $L = \infty$ limit provide the missing link for the qD-coupling of $L^d = 0^d$ states, and thus 'complete' the set of $SU(6) \times O(3)$ qD-couplings of baryons. To see this explicitly we reproduce the full widths of $J = L \pm 1/2$ N^d -states as a sum

of contributions from the $qD_S(S)$ and $qD_\mu(A)$ modes as follows :

$$\Gamma_S^{d(L+1/2)}(N) = (3/4) [L/(L+1)] \alpha_S Q \sqrt{2} (\omega_1 - m_1) / M ;$$

$$\Gamma_A^{d(L+1/2)}(N) = 3 \Gamma_S^{d(L+1/2)}(N) \quad (42)$$

$$\begin{aligned} \Gamma_S^{d(L-1/2)}(N) &= (3/4) [(L^2 - L + 4)/(L+1)^2] \alpha_S Q \sqrt{2} (\omega_1 + m_1) / M ; \\ \Gamma_A^{d(L-1/2)}(N) &= (9/4) [(L^2 - L + 2)/(L+1)^2] \alpha_S Q \sqrt{2} (\omega_1 + m_1) / M \end{aligned} \quad (43)$$

where

$$\alpha_S = g_S^2 / 4\pi ; \quad \omega_1^2 = m_1^2 + Q^2 . \quad (44)$$

Now taking the limit $L = \infty$, after averaging over the $L \pm 1/2$ contributions, we obtain a rather simple formula (with $\sqrt{2} = 1$) :

$$\Gamma^d("L = 0") = 3\alpha_S Q \omega_1 / M \quad (45)$$

which, e.g., gives ~~337~~ MeV for the N(1440) width (vs. 350 ± 100 expt) with no adjustment whatsoever; (see discussion in Sec.4.2 below).

In Appendix B are listed the principal formulae for the baryon total widths, including the flavour factors {FL}.

4. Results and Discussion :

We are now in a position to present the numerical results for comparison with data [1] in respect of both types of hadrons in several flavour sectors, as well as their interpretation w.r.t. the theoretical premises outlined in Sec.1. To that end we quickly recapitulate the basic inputs, viz., the masses of the 'quasi-free' quarks and diquarks on the lines explained in Sec.1.2 already :

their values (instead of being entirely free parameters) were determined from a chirally symmetric BSE model [5] attuned to [7] spectroscopy [1] and in conformity with Politzer additivity [10]. For the diquark masses too, the general quark-model additivity was employed to determine their various combinations. The final inputs (in MeV) (with $q = u/d$) are :

$$m_q = 176, m_s = 326 ; m_D(qq) = \underline{352} ; m_D(qs) = \underline{502} ; m_D(ss) = \underline{652}. \quad (46)$$

As to the overall coupling constant g_s , a tentative identification was suggested (see Sec.1.3) in terms of the standard QCD value α_s merely to avoid an unnecessary proliferation of parameters without good reason. With 3 flavours which fits the spectra [1]:

$$g_s^2/4\pi = \alpha_s(M^2) = 2\pi / [9 \ln (M/\Lambda_{\text{QCD}})] ; \Lambda_{\text{QCD}} = 200 \text{ MeV [7]}. \quad (47)$$

4.1 Meson Sector :

In Sec.2 we have considered the coupling of mesons to both $q\bar{q}$ and DD^* channels. The $q\bar{q}$ coupling information is compactly contained in eqs. (3,14). For the DD^* coupling, the classification scheme for the available channels is outlined in Sec.2.3 together with postulated selection rules for natural parity states of "good" G-parity. (See Appendix A for calculational details on DD^* couplings, in accordance with Partial Symmetry [12,13]; and Appendix B plus Table 1 for organization of the width formulae for various cases, including the effect of different flavour species) .

The results on meson widths are given in Tables 2,3.4 for

non-strange, strange and hidden-strange cases respectively, in the standard PDG [1] notation for the various states involved. For the isovector (ρ , π and a types) and isospinor (K and K^* types) a good pattern of accord of the data with "predictions" is seen, subject (of course) to the variations in the data themselves. In this respect there has been no special selection rule for the K , K^* , π series, but there seems to be a distinct bearing of our postulated 'selection rule' for the ρ, a series on the accord with data for the available states (the ground state ρ ($J=1$) shortfall due to phase space limitations notwithstanding). Note that the same selection rule is also responsible for the large width of the a_1 and the small width of f_1 . For the kaon ($I=1/2$) cases Table 4 gives "pure" assignments (in parallel with $I=1$) instead of 50-50 admixtures suggested in [1].

A more sensitive test of the coupling strengths comes from the $q\bar{q}$ and associated DD^* couplings of the isoscalar (ω , f) and (ϕ , f') towers wherein their full strengths $g_S\{1/3; 1/3+S\}$ for $q\bar{q}$ and $g_S\{2/3; 2/3+S\}$ for DD^* are determined by their respective baryonic charges and hypercharges [11]. Note that these values are quite distinct from those of the $SU(6) \otimes O(3)$ quark model [25], (being almost a third of the latter). For the L-excited ω and ϕ these "Sakurai" strengths, in association with the postulated selection rules for their DD^* couplings, seem to be just right for reproducing the pattern of the data, thus ipso-facto showing a significant role for these "gauge-couplings"

with strength g_s ; see eq.(5). Unfortunately these Sakurai [11] strengths do not seem to extend all the way to the ground states ($L=0$) of ω and ϕ : for these two states we have had to take only half the strength ($g_s/2$), which amounts to Schwinger's e.m. substitution rule [12] with the correspondence " e " $\leftrightarrow$ " g_s ". The price of extending the Sakurai strengths all the way to $L=0$ would have been to give four times the observed ω and ϕ widths.

We should now comment on the other more general aspects of our coupling scheme. Thus it is seen from Tables 2-4 that the $q\bar{q}$ mode gives a rather modest contribution to the total widths of most mesons , being at most a third for the ρ, K^* series (the largest contributors) and almost completely negligible for the K, b, π series. The dominant contributions seem to come collectively from the DD^* channels which, though individually quite modest each, nevertheless make up through the existence of several distinct combinations of them for each hadron (see Sec.2.3 and Appendix B). In this connection we also note the conjectured importance of four-quark states even for relatively low excitations [21].

On the negative side, our results for low J states , especially $J = L - 1$, are consistently on the "low" side . For though we seem to 'understand' the widths of f_0 and a_0 , (presumably an artifact of the phase space), the predictions are consistently smaller than the data for the higher excitations [$\omega'(1600)$, $f_0(1590)$ and $\rho'(1700)$]. On the other hand our predictions on

radially excited states like $\omega_R(1420)$ and $\rho_R(1450)$ - see page 1320 of [1] - are somewhat on the high side, thus suggesting the possibility of mixing, e.g. $\omega(1420)$ with $\omega(1600)$, etc.

The problem of low widths of $J = L-1$ states may be traced back to the requirement of gauge invariance in their coupling structures vis-a-vis those of $J = L+1$ and L of natural parity : For the latter states our method of construction [see Appendix A] yielded gauge-invariant couplings in a "natural" manner but for the former ($J=L-1$), a constraint has had to be imposed by keeping only the transverse parts of the D_μ, D_μ^* fields, thus resulting in a significant reduction in the effective couplings of the S^J - mesons to the diquark fields. For a gauge invariant construction an alternative possibility would have been via standard "gauge substitutions" like $\partial_\mu \rightarrow \partial_\mu - i g_{\text{eff}} S_\mu^J$ in the 'free' D and D^* fields, but such methods would not have generated the 'mixed' couplings like S_{Jm} (see table 1) which have proved important contributors to the overall 'good' pattern of results. (It would thus seem that the Partial Symmetric method [12] of generating DD^* couplings from the $q\bar{q}$ meson -matrix, vide Appendix A, is more microscopic than the direct gauge substitutions on the corresponding 'free' D, D^* fields).

The last two aspects of our coupling scheme concern the roles of the "velocity form factor" v^L [14] on the one hand , and the (tentative) identification of the overall coupling strength g_s

with its QCD value on the other . Although it is difficult to pin-point the sensitivity of their individual roles in producing the overall agreement, their variations over large masses and L-values have nevertheless 'registered' in the process. [A simple way to "check" this would be to "suppress" these factors and find a significant distortion in the agreement !]

4.2 Baryon Sector :

We now turn to the other side of the coin, viz., the qD coupling of baryons which was formulated in Sec.3 on the basis of an implied 'duality' between a diquark and an anti-quark (both color 3^*), except for ignoring the $q\bar{q}$ Fock component (dual to DD^*) in the baryon. The $SU(6) \otimes O(3)$ couplings of baryons , eqs.(38-39), which were inferred from their qD contents eqs.(26-31), are again characterized by the same overall coupling constant g_s as well as the same input quark masses, viz., eqs.(46-47) above, and the results are summarized in tables 5 & 6 for non-strange (N, Δ) and strange (Λ, Σ) baryons respectively.

Because of the highly concise structure of the baryon qD coupling facilitated by the absence of the higher Fock components, the results in tables 5 & 6 are quite straightforward to interpret in terms of the standard assignments. In this respect our coupling scheme does distinguish among the $56^{q,d}$ and $70^{q,d}$ types , and the same have been indicated in tables 5 & 6

which by and large conform to the PDG [1] list (p 1322).

We do find good overall accord with the data in both tables without any additional assumptions and/or parametrizations. This is separately true of the $q\bar{D}$ couplings of L^d and L^q states, especially the trends for higher L -excitations. In particular, for those cases where the data are quite precise, e.g., $\Lambda(1405)$, $\Lambda(1520)$, $N(1520)$ and $\Sigma^*(1385)$, the predicted values seem to follow the data quite closely. (For the 'strange' baryons, we have fully incorporated the effect of permutation symmetry by taking the s -quark as a component of the diquark $2/3$ of the time, and the u/d quarks to make up the latter only $1/3$ of the time).

The only outstanding area of discord concerns the N -states of 70^q whose widths seem to be 2-3 times the observed values. These cases may be a pointer to the possible role of some 'selection rules' (similar to the meson cases of 'good' G -parity) designed to reduce the weightage of the $I = 1$ diquark. Such possibilities in turn depend on the (hitherto neglected) higher Fock components such as $q\bar{D}q\bar{q}$ since there is practically no freedom in our minimal coupling scheme which is quite rigidly defined. (We have not attempted this exercise here).

Now to some general remarks concerning the coupling structure. In the language of 'multipole fields' [24], the 'electric' couplings

are responsible for the decays of 56^q and 70^q states, while the 'magnetic' couplings account for those of 56^d and 70^d . As to the two quantitative ingredients, the coupling constant α_s and the velocity factor v^2 , their role has been very much the same as in the meson case : they have contributed significantly to the overall agreement . The only cases that could not be included in the scheme have been qD couplings of $L^d = 0^d$ states which apparently vanish for $L=0$ (see eq.(39)). This problem was sought to be circumvented by considering, not the $L = 0$ limit, but rather the opposite limit $L=\infty$, as a possible candidate to apply to such cases as N(1440) or N(1710), as if to suggest that the route to $L = 0$ lies via $L = \infty$! We could not offer any deep reason for this procedure, except for its semi- intuitive appeal and the natural agreement of its predictions on N(1440) and N(1710) with the data, without having to bring in a fresh coupling constant .

We should also like to comment on the case of a rather special baryon, $\Delta(1930)D_{35}$, from the viewpoint of understanding its status as a possible 56 state of natural odd - parity. We had recently examined the mass location of this state [26], but found the theoretical value too high for the observed location, in apparent agreement with a much earlier finding [27]. In the present analysis we have considered two parallel assignments , viz., $[70^d, 3^-]$ and $[56^q, 1^-]$ as the simplest possibilities within our scheme. The results for the two assignments are 107 MeV and 292 MeV respectively , versus 250-450 [1], thus showing a clear

preference for the 56^- assignment. In view of the often greater sensitivity of resonance widths to their spectroscopic assignments, than of their mass locations, we are inclined to regard this result as a possible signature for a 56^- assignment. The other two candidates for more rigid 56^- assignments [26], viz. $G_{39}(2400)$ and $I_{3,13}(2750)$ seem to give qualitatively correct pictures for the observed widths.

4.3 Γ/M Ratios for Hadrons :

After this fairly convincing set of agreements in several distinct sectors of strong interactions, it is natural to ask what has been learnt from this (rather elaborate) exercise, and in particular if the result of this investigation can be put in a nutshell with possible bearing on its further ramifications. One such item which brings out the close interrelation between Γ and M on the basis of the present investigation is the ratio Γ/M for which a rather striking formula of the form $n/2^n$ has been found by RS [3], over the entire spectrum of elementary particle decays including the strong interaction sector. It should therefore be of considerable interest to ask how our Γ/M ratio, derived from entirely orthodox theoretical premises, would look like in the strong interaction domain. Now from the various width formulae derived in the text, it is easy to check that for hadron masses large compared to the quark and diquark masses ($Q \approx M/2$), the ratio $R \equiv \Gamma/M$ assumes a very simple form,

separately for each hadronic species. For example for some principal meson series (π, ρ, K, K^*) one finds

$$R_{\pi} = 2R_K = \alpha_s/2 ; R_{\rho} = 2R_{K^*} = (5\alpha_s/4) [J+1/2J+1] ; \quad (48)$$

where for 3 flavours (see eq.(47))

$$\alpha_s^{-1} = (9/2\pi) \ln (M/\Lambda_{\text{QCD}}) \quad (49)$$

Similarly for the baryon series N^d ($J=L+1/2$) one finds

$$R(N^d) = (3\alpha_s/4) [(J - 1/2) / (J + 1/2)] . \quad (50)$$

On the other hand for the baryon series Δ^q ($J=L+3/2$) the ratio increases with M , and also depends on the diquark mass m_D :

$$R(\Delta^q) = (3\alpha_s/16) [1/2 - 1/(4J)] M^2/(4 m_D^2) . \quad (51)$$

Similar formulae hold for other flavour species.

An inspection of these formulae suggests that to the extent that the running coupling constant α_s involves the ratio M/Λ_{QCD} , the 'discreteness' or otherwise of the ratio R is directly reflected in the corresponding behaviour of the hadron mass M . The latter in turn has long been known to exhibit the standard Regge slope behaviour which is also upheld by the quark model [2,25]. Thus our result is in general conformity with the RS [3] finding of discreteness in the ratio R in the strong interaction sector, inasmuch as the same is made to depend on that of M/Λ_{QCD} , but a more quantitative connection with "n" (as distinct from a mere numerical evaluation) is not immediately in sight.

5. Summary and Conclusion :

Looking back on the results on both types of hadronic decays, it is perhaps worth noting the strong similarity of their respective constitutions which is best described by an underlying duality principle and has been at the root of this unified approach to the study of their decay characteristics. Indeed the results described in the foregoing with a common set of postulates and inputs would seem amply to justify such a "duality" interpretation. We must also keep in mind that the entire exercise has been centred around the basic ansatz of a short 'breathing mode' for the temporary dissociation of a hadron into a pair of quasi-free constituents [5] before eventual hadronization (with certainty). Therefore the outcome must be taken to have a strong causal connection of this basic hypothesis in association with our "minimal" dynamical scheme of Yukawa couplings through which it has been revealed . In this respect the rather impressive agreement of the calculated widths on a whole range of hadronic sectors simultaneously , with just two inputs, eq.(46-47), would seem to provide a strong degree of support to the hypothesis itself, since the basic premises [11,12] within which the dynamical scheme has been presented are already time-tested. Nevertheless certain technical aspects of the scheme , viz., multiple derivative couplings [19,20] and velocity form factors [14,15] , together with the QCD "running coupling constant" have all played their silent roles in achieving the

goal of providing a semblance of order to an otherwise "diffuse" collection of data in the strong interaction sector. The latter is summarized in eqs.(48-51) in the form of $R \equiv \Gamma/M$ for the different hadronic species, which is suggestive of discreteness, an aspect which is best left for further studies.

One of us (ANM) is indebted to Prof.Raja Ramanna for 'inspiration' to look at the problem of hadronic widths, in the context of his recent theory [3]. We are also grateful to him for the warm hospitality of the Institute.

Appendix A :

The meson-matrix, eq.(1) of text [11] represents the couplings of the positive parity components of the corresponding mesons, considered as radiation quanta, to a primary $q\bar{q}$ pair. These primary couplings can be folded into various composite structures (baryons, mesons etc.) by standard quark model techniques, with subsequent Lorentz boosting. For the ground state meson couplings the results were first described in [18]. A subsequent generalization to a much larger variety of meson and baryon couplings, especially an extension to their L-excited states, with inclusion of quark recoil and convective current terms, eq.(3) of

text, was described in a more systematic way in a longer review [13]. In this appendix we summarize the main aspects of these derivations [13] as adapted to the couplings of the radiation quanta -eq.(1)- and their L-excited counterparts, to the DD^* pairs as quark composites.

qqbar-couplings of mesons :

First, for the $q\bar{q}$ couplings themselves, the Lorentz boosting is almost self-evident with the replacements

$$\vec{\sigma} \cdot \vec{k} \pi^a \rightarrow 2i\gamma_5 \gamma \cdot \vec{Q} \pi^a \quad ; \quad V_0^a \rightarrow i\gamma_\mu V_\mu^a, \text{ etc.} \quad (\text{A.1})$$

all sandwiched between q and $\bar{q}$ 4-spinors. The term $\vec{\sigma} \cdot \vec{A}$ in eq.(1) for the positive parity axial counterpart (A_i^a) of V_i^a , becomes $i\gamma \cdot A \gamma_5$, with equal strength as appropriate for chiral symmetry. The magnetic V-term in eq.(1) can be similarly boosted, but care must be taken to avoid double counting at the quark level, since the (normal) magnetic moment of the quark is already contained in the $i\gamma \cdot V$ term. [However, this term may contribute (without duplication) to the coupling of mesons to quark composites such as diquarks (see below)]. The flavour structure which is fully defined in terms of λ -matrices in eq.(1), may be carried over without change to the boosted form, leading finally to eq.(4) of text for ($L=0$) meson- $q\bar{q}$ couplings.

For the $q\bar{q}$ of L-excited mesons, the same logic applies without change except that for each L there are now four classes of mesons

listed in the text as spin-singlet B^L -states and spin-triplet V^J, A^J, S^J states of $J = L+1, L, L-1$ respectively. As already noted in sec. 2.1, these states are best expressed in their "Q-contracted" forms. Q_μ is the relative 4-momentum of the outgoing $q\bar{q}$ pair of 4-momenta $p_{1\mu}, p_{2\mu}$ and masses m_1, m_2 so that

$$p_{1,2}^\mu = \hat{m}_{1,2} P_\mu \pm Q_\mu ; Q \cdot P = 0 ; 2 \hat{m}_{1,2}^2 = 1 \pm (m_1^2 - m_2^2) / M^2 \quad (A.2)$$

The rest of the derivation including the role of the velocity form factors [13,14] is as given in eqs.(5-14) of text. After the velocity factor is extracted, Q_μ appears in the tensorial structures as a unit vector $\hat{Q}_\mu$ in most of the formulae.

Meson Couplings to DD^* pairs :

We now sketch briefly the DD^* couplings of both ground and L-excited mesons in the light of the selection rules outlined in sec. 2.2 of text.

The coupling scheme for DD^* pairs follows the 'classical quark model' pattern [18] of coupling of 'radiation quanta' to mesons and baryons [13] via the additivity principle, the only difference being that a 'diquark' now replaces the $q\bar{q}$ or qqq hadrons. To that end the basic meson matrix (1) for quark couplings may be folded into the D^*D structures in the limit of low momentum transfers, (with the mesons themselves playing the role of radiation quanta). The subsequent 'boosting' can be done as in the analogous $q\bar{q}$ cases of (3) and (14). We illustrate this

procedure for some typical cases, viz., the $DD^*- \rho$ and DD^*-b_1 couplings. In general for a DD^*-B^L coupling, taking the relevant term from (1), the standard quark model [18,13] in the low momentum limit yields the spin-isospin trace (here the λ -matrices are replaced by isospin τ -matrices)

$$[\text{Two}] \xi(L) "VW" g_S/2 M_L \text{Tr} [\vec{\sigma} \cdot \vec{D}^{*a} 2\vec{\sigma} \cdot \vec{D}^{Lc} \vec{\sigma} \cdot \vec{D}^b/2] \times \text{Tr} [\tau_a \tau_b \tau_c/2] \quad (\text{A.3})$$

Here the normalized diquark states (both axial vectors) are represented in spin-isospin space as $\vec{\sigma} \cdot \vec{D}^{*a} \tau_a / 2$ and $\vec{\sigma} \cdot \vec{D}^b \tau_b / 2$; the factor of 'two' in front refers to the summation over two quarks in the diquark; and the factor "VW" stands for the normalization [18] needed to provide the 'correct' dimension to a 3-boson coupling (on par with boson-fermion coupling). The question is which 'particle' must be subjected to the "VW-effect" since this factor is supposed to 'bring up' the particle from the 'state' to the 'field' status. In the present case we have two such particles (the diquarks) to upgrade to the field status, and this must be done symmetrically. This last requirement, following ref. [18] in the low momentum limit of these particles, yields the fairly unique value $\sqrt{2m_1 2m_2}$. (Other possibilities also exist, see ref. [13], for the choice of the radiation quantum, in which case the VW-factor would be different.) The Lorentz-boosted form of eq.(A.3) after due simplification is given in eq.(17) of text. Note that a B^L does not couple to $D_S^* D_S$ pairs, but does to $D_\mu^* D_S$, and the analogous coupling structure may be worked out as in (A.3) with D_S and D_S^a states represented in the 2x2 forms $D_S \times I/2$ and

$D_s^a \tau_a / 2$ respectively. Its Lorentz-boosted form is given in eq.(18) of text. A multiplicative factor W_L with

$$W_L = \sqrt{4m_1 m_2} / M_L \quad (A.4)$$

is associated with all the Lorentz boosted DD^* couplings listed in the text. The A^J couplings which also come from two channels, viz., $D_\mu^* D_\nu$ and $D_s^* D_\mu$ can be taken over almost verbatim from the above derivation for the B^L -series through the replacement $\sigma \cdot Q \rightarrow \sigma \cdot A^J$ in the meson matrix, eq.(1), with resulting coupling forms $\tilde{a}_J$ and $\tilde{a}_{Jm}$ respectively. The result for $\tilde{a}_J$ is shown in eq.(24) of sec. 2.2. The $\tilde{a}_{Jm}$ coupling on the other hand comes out proportional to $D_s^*(p_2) A_\mu^J(Q) D_\mu(p_1) + \text{h.c.}$ which unfortunately is not yet gauge invariant. This defect of the derivation can however be remedied in the standard mannner $\delta_{\mu\nu} \rightarrow \delta_{\mu\nu} - \hat{Q}_\mu \hat{Q}_\nu$, and the final result is given in eq.(25) of sec.2.2.

Consider next the case of V^J -mesons which can couple to all the combinations $D_\mu D_\nu^*$, $D_s D_s^*$, $D_\mu D_s^*$ (and h.c.). The coupling to $D_\mu D_\nu^*$ can be worked out exactly as in (A.3) via the magnetic V-term, but now the contribution comes only from the recoil term of $P_i = M p_i^{(1,2)} / \omega_{(1,2)}$ (see eq.(4) of text) in the rest frame $P_i = 0$ of the hadron. Similarly the $D_s D_s^*$ coupling to a V-meson may be obtained from the 'charge' term V_0 of eq.(1) via the 'convective' [2] contribution, eq(4). For the L-excited V^J mesons too, the procedure is the same as already defined in the text for

their $q\bar{q}$ couplings, viz., the CG reduction eq.(10) can be folded into the corresponding coupling types merely by multiplying with the relevant factors $\eta(L)$ and $\zeta(L)$ associated with the lower ranking A^J ($J = L$) and S^J ($J = L-1$) terms respectively in the CG expansion, eq.(9).

For the S^J -terms, there appear fresh coupling structures $D_\mu^* S^J D_\nu$, $D_S^* S^J D_\nu$ and $D_S^* S^J D_S$, not yet encountered in the above derivation. Being a lower ranking meson, S^J couplings to these diquark pairs come from two distinct pieces of the meson matrix (1) : i) the magnetic V-term (via recoil effect) and ii) the charge V-term (via convective effect). However a difference of these couplings from the corresponding V-meson couplings noted above arises from the fact that the "field" V_μ is now replaced by the generalized vector

$$\tilde{V}_\mu \equiv \zeta(L) [2 Q_\mu S^J(Q) + (L-1) Q^2 S_\mu^J(Q)], \quad (A.5)$$

as may be immediately checked from the S^J term of the CG expansion, eq.(10) of text. The first term of (A.5) does not couple to the magnetic term of the meson-matrix (1); only the second one does and gives rise, via the recoil contribution, to a coupling of the form ρ_J , eq.(19). However the first term ($\sim Q_\mu$) can contribute, via the convective part [2] of the 'charge' term, eq.(4), a new coupling proportional to $D_\mu^*(p_2) D_\mu(p_1)$. Since on the other hand this coupling form lacks gauge invariance, we may exploit the inherent ambiguity in these high

spin $(D_\mu^* D_\nu)$ couplings of a low ranking S^J to replace $D^*.D$ by $D^*.D - D^*.Q D.Q/Q^2$ to make it explicitly gauge-invariant (analogously to the $\tilde{a}_{Jm}$ coupling noted above). The final structure (with all factors shown explicitly) is listed in eq.(23) of text . For a complete description of the flavour-cum-Lorentz structures of all the DD^* couplings see sec.2.3 of text in conjunction with Appendix B (next).

Appendix B :

We collect here the necessary formulae for the decay widths of mesons and baryons on the basis of their coupling structures described in the text.

qqbar modes of Mesons :

For the $q\bar{q}$ modes the relevant formulae for the (B^J, V^J, A^J, S^J) types are (c.f. [5]) :

$$\Gamma_{q\bar{q}}^-(B^J) = \alpha_s [FL] Q \sqrt{2} \left(1 - \frac{\delta m^2}{M^2} \right) \frac{m_{12}^2}{M^2} ; \{m_{12}=m_1+m_2, \delta m = m_1-m_2\}$$

$$\Gamma_{q\bar{q}}^-(V^J) = \alpha_s [FL] Q \sqrt{2} \left(1 - \frac{\delta m^2}{M^2} - \frac{4J}{2J+1} \frac{Q^2}{M^2} \right) ; J = L + 1$$

$$\Gamma_{q\bar{q}}^-(S^J) = \alpha_s [FL] Q \sqrt{2} \left[\left(1 - \frac{\delta m^2}{M^2} \right) (J(2J+1) + 4(J+1)) / (J+2)^2 - \frac{4Q^2}{M^2} \right] \frac{L}{2L-1} ; J = L - 1$$

$$\Gamma_{q\bar{q}}^-(A^J) = \alpha_s [FL] Q \sqrt{2} \left(1 - \frac{m_{12}^2}{M^2} - \frac{2J-1}{2J+1} \frac{4Q^2}{M^2} \right) \frac{J}{J+1} ; J = L$$

(B.1)

α_s is the QCD coupling constant given by eqs.(44,47) of the text
 [FL] is a 'flavour factor' whose values for the different meson
 types are listed below :

$$\begin{array}{ccccccc}
 & \pi & K & ; & (\rho/a_1) & K^* & (\omega / f) & (\phi / f') \\
 [FL]_0 = & \frac{1}{2} & \frac{1}{4} & ; & \frac{1}{2} & \frac{1}{4} & \frac{1}{18} ; 0 & \frac{1}{9} ; 0 \\
 [FL]_L = & \frac{1}{2} & \frac{1}{4} & ; & \frac{1}{2} & \frac{1}{4} & \frac{2}{9} ; \frac{2}{9} & \frac{4}{9} ; \frac{4}{9}
 \end{array}
 \tag{B.2}$$

These two lines refer to the $L=0$ and L -excited mesons
 respectively, as read from the λ -matrices defined in eqs.(2) and
 (5) of the text.

DD^* modes of Mesons :

In the notation of sec.2.3 for the different types of meson DD^*
 couplings [see eqs.(17-25) of text], the general formula for their
 DD^* widths may be worked out in a straightforward manner using the
 formulae (11-13) for the meson projection operators and the
 corresponding diquark projection operators ($i=1,2$)

$$\Theta_{\mu\nu}(p_i) = \delta_{\mu\nu} + p_{i\mu} p_{i\nu} / m_i^2. \tag{B.3}$$

Omitting the calculational details for brevity the final result
 is organised as under.

$$\Gamma_{DD^*} = \{FL\} C(J) X(L) \tag{B.4}$$

where

$$X(L) = \alpha_s Q W_L^2 [4Q^2/M^2]^{(L+1)} ; W_L^2 = 4 m_1 m_2 / M^2 \tag{B.5}$$

The quantity $C(J)$ for the four different towers may be compactly
 expressed as under the following table :

	$\{J\}$	$\{J_m\}$	$\{J_s\}$
B^J	1	$(1+\theta)$	0
V^J	$\frac{J+1}{2J+1} (1+\theta)$	$\frac{J+1}{2J+1}$	$\frac{J/2}{2J+1}$
A^J	$(1 + \frac{J+1}{2J+1} \cdot \theta) \frac{M^2}{4Q^2}$	$\frac{J+1}{2J+1} \frac{M^2}{2Q^2}$	0
S^J	3 (J+2)	0	$\frac{J+2}{2}$

where

$$2\theta = (Q/m_1)^2 + (Q/m_2)^2$$

The factor $\{FL\}$ in eq.(B.3) which distinguishes the different flavour modes of DD^* dissociation is arranged in table 1.

qD modes of Baryons :

The formulae for the baryonic widths may be worked out in close analogy to the meson case , but the details are again omitted for brevity. Note that the Ψ^{Jq} states in our scheme couple only to qD_μ pairs [not qD_s]. The general formula for all cases may be written as a product of three factors :

$$\Gamma_{A,S}^J = Y(L) \{FL\} H(L) \quad ; \quad (A = \text{axial} , S = \text{scalar}) \quad (B.6)$$

$$Y(L) = (3/4) \alpha_s Q \sqrt{2L} \quad ; \quad G_\pm = (\omega_1 \pm m_1) / M \quad (B.7)$$

$\{FL\}$ is a flavour factor with values $2(N), 1(\Delta), 1/2(\Lambda), 1/2(\Sigma)$ for the respective quartet (q) states. For the doublet (d) states the corresponding values for A- and S- couplings of **56** and **70** baryons are :

56^d : N (3/2, 1/2) ; Δ (0,0) ; Λ (0,1/2) ; Σ (1/2,0)

70^d : N (1,1) ; Δ (1/2,1/2) ; Λ (1/4,1/4) ; Σ (1/4,1/4)

For the (d) baryons the formulae are explicitly given in eqs.(42) and (43) of text for the N-case. For the (q) baryons the function $H(L)$ is given as follows :

$$\begin{aligned} J=L+3/2 & \quad G_- [1 + (L+1)\theta/(2L+3)] ; \\ J=L+1/2 & \quad G_+ [2L^2 + 3L + 3 + (L+1)^2\theta]/[(L+1)(2L+3)] ; \\ J=L-1/2 & \quad G_- L [2L^2 + L + 1 + (L+1)^2\theta]/[(L+1)^2(2L-1)] ; \\ J=L-3/2 & \quad G_+ L [1 + \theta]/(2L-1). \end{aligned} \quad (B.8)$$

Finally for the qD couplings of Λ and Σ states it is necessary to lift the restriction of #1 for the s-quark as the basis state, vide eqs.(29-31), so as to restore permutation symmetry. This is done simply by taking the diquark to be made up of qq-quarks (q=u/d) one-third of the time, and of the sq combination two-thirds of the time. The corresponding diquark masses are 352 and 502 MeV respectively

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Table 1

Two-way classification of {FL} according to Lorentz structures $(V_J, V_{Jm}, V_{Js} ; B_J, E_{Jm} ; A_J, A_{Jm} ; S_J, S_{Jm})$ distributed among the different flavour combinations considered in the DD^* pairs in terms of their non-strange $(q, \bar{q})$ and strange $(s, \bar{s})$ contents. For the natural parity states of "good" G-parity the odd and even L-excited states are shown separately in accordance with the selection rules outlined in sec.2.3.

		$q\bar{q}^{2-2}$	$qs\bar{q}^{-2}$	$s\bar{s}^{2-2}$
$(\rho_J; \rho_{Jm}; \rho_{Js})$	$(J=1,3,5..)$	$(2;1;0)$		
$(a_J; a_{Jm}; a_{Js})$	$(J=0,2,4..)$	$(0;1;2)$		
$(\tilde{a}_J; \tilde{a}_{Jm})$	$(J=1,3,5..)$	$(2;1)$		
$(\omega; \omega_s)$	$(J=1 \text{ only})$	$(\frac{1}{3}; \frac{1}{9})$		
$(\omega_J; \omega_{Js})$	$(J=3,5..)$	$(\frac{4}{3}; \frac{4}{9})$		
$(f_J; f_{Js})$	$(J=0,2,4..)$	$(\frac{16}{9}; \frac{16}{9})$		
$(b_L; b_{Lm})$		$(2;1)$		
$(K_J; K_{Jm}; K_{Js})$	$(\text{all } J's)$		$(1; 1/2; 0)$	
$(\phi_J; \phi_{Jm}; \phi_{Js})$	$(J=3,5..)$			$(\frac{16}{9}; \frac{16}{9}; \frac{16}{9})$
$(f'_J; f'_{Jm}; f'_{Js})$	$(J=2,4..)$			$(\frac{16}{9}; \frac{16}{9}; \frac{16}{9})$

TABLE 2: Total Widths (in MeV) of non-strange mesons with contributions from different channels. For notations see text and Appendix B. (subscript R =radial excitation; prime denotes $J=L-1$ states.)

State	$\Gamma(\text{expt.})$	$\Gamma(q\bar{q})$	Γ_J	Γ_{Jm}	Γ_{Js}	$\Gamma(\text{tot})$
$\rho(770)$	151 ± 12	60.9	16.4	6.9	0	84.2
$\omega(782)$	8.43 ± 1.1	6.80	2.73	0	0.07	9.60
$f_0(980)$	40-400	5.00	36.2	0	0	41.2
$a_0(980)$	50-300	11.3	54.3	0	0	65.6
$h_1(1170)$	360 ± 40	9.0	108.8	300.6	0	418.4
$b_1(1235)$	142 ± 8	8.5	57.6	88.6	0	154.7
$a_1(1260)$	~ 400	35.2	306.1	56.7	0	398
$f_2(1270)$	185 ± 20	29.8	151	0	15.5	196.3
$f_1(1285)$	24 ± 3	15.9	12.0	0	0	27.9
$\pi(1300)$	200-600	8.6	84.5	144	0	237.1
$a_2(1320)$	107 ± 5	68.7	0	18	12.1	98.8
$\omega_R(1420)$	174 ± 60	37.0	142.6	0	11.9	191.5
$\rho_R(1450)$	310 ± 60	85.0	247.7	26.8	6.7	366.2
$\omega'(1600)$	280 ± 24	12.2	34.9	0	1.2	48.3
$\omega_3(1670)$	173 ± 11	32.1	101.3	0	17.4	150.8
$\pi_2(1670)$	240 ± 15	5.1	49.2	138.5	0	193
$\rho_3(1690)$	215 ± 20	73	162.1	14.1	0	249.2
$\rho'(1700)$	235 ± 50	26.4	64.4	7.4	1.8	100
$f_4(2050)$	208 ± 13	35.0	186.0	0	9.2	230.2

TABLE 3 : Total widths of (ϕ, f') mesons; for notations see table 2

State	$\Gamma(\text{expt.})$	$\Gamma(q\bar{q})$	Γ_J	Γ_{Jm}	Γ_{Js}	$\Gamma(\text{tot})$
$\phi(1020)$	4.43 ± 0.06	3.75	0	0	0	3.75
$f_1(1510)$	35 ± 15	46.0	27.2	29.5	0	102.7
$f_2'(1525)$	76 ± 10	58.0	10.5	7.7	0	76.2
$\phi'(1680)$	150 ± 50	26.1	8.7	1.7	0.4	36.9
$\phi_R(1680)$	150 ± 50	80.9	81.9	49.3	12.3	224.4
$\phi_3(1850)$	87^{+28}_{-23}	57.9	21.3	7.9	2.9	90
$f_2'(2300)$	149 ± 40	29.8	30.4	7.4	2.5	70.1
$f_{2R}(2340)$	319^{+80}_{-70}	89.5	139.8	43.3	14.5	287.1

TABLE 4 : Total widths of strange mesons; for notation see table 2

State	$\Gamma(\text{expt.})$	$\Gamma(q\bar{q})$	Γ_J	Γ_{Jm}	Γ_{Js}	$\Gamma(\text{tot})$
$K^*(892)$	50	31.9	6.2	2.8	1.4	42.3
$K_1(1270)$	90 ± 20	8.2	4	51.4	0	63.6
$K_1(1400)$	174 ± 13	16.1	98.0	59.1	0	173.2
$K_2^*(1430)$	109 ± 5	37.9	49.2	16.6	5.5	109.2
$K_R^*(1680)$	323 ± 110	45	173.9	42.2	10.5	271.6
$K_2(1770)$	186 ± 14	4.5	24.7	113.2	0	142.4
$K_3^*(1780)$	164 ± 17	33.8	130	14	10.4	188.2
$K_4^*(2045)$	198 ± 30	35.1	146.5	11.9	9.5	203

TABLE 5 : N, Δ widths (in MeV) with axial (A) and scalar (S) diquark contributions shown separately.

State	$SU(6); L^P$	$\Gamma(\text{expt.})$	Γ_A	Γ_S	$\Gamma(\text{tot})$
$\Delta(1232)$	$(56, 0^+)^q$	120 $\pm$ 5	93.2	0	93.2
$P_{11}(1440)$	$(56, 0^+)^d$	350 $\pm$ 100	252.6	84.2	336.9
$D_{13}(1520)$	$(70, 1^-)^d$	110-135	57.9	57.9	115.8
$S_{11}(1535)$	$(70, 1^-)^d$	100-250	95.3	190.4	285.7
$P_{33}(1600)$	$(56, 0^+)^q$	250-450	176.8	0	176.8
$S_{31}(1620)$	$(70, 1^-)^d$	120-180	49.7	99.4	149.1
$S_{11}(1650)$	$(70, 1^-)^q$	145-190	756.5	0	756.5
$D_{15}(1675)$	$(70, 1^-)^q$	140-180	396	0	396
$F_{15}(1680)$	$(56, 2^+)^d$	120-140	118.5	39.5	158
$D_{13}(1700)$	$(70, 1^-)^q$	50-150	596	0	596
$D_{33}(1700)$	$(70, 1^-)^d$	200-400	66.9	66.9	133.8
$P_{13}(1720)$	$(56, 2^+)^d$	100-200	123.8	61.9	185.7
$F_{35}(1905)$	$(56, 2^+)^q$	280-440	378.5	0	378.5
$P_{31}(1910)$	$(70, 0^+)^d$	190-270	103.3	103.3	206.6
$P_{33}(1920)$	$(56, 2^+)^q$	150-300	258	0	258
$D_{35}(1930)$	$(56, 1^-)^q$	250-450	291.7	0	291.7
$D_{35}(1930)$	$(70, 3^-)^d$	250-450	47.7	59.6	107.3
$F_{37}(1950)$	$(56, 2^+)^q$	290-350	290.6	0	290.6
$G_{17}(2190)$	$(70, 3^-)^d$	350-550	118	118	236

contd.

TABLE 5 (in continuation)

State	$SU(6); L^P$	$\Gamma(\text{expt.})$	Γ_A	Γ_S	$\Gamma(\text{tot})$
$H_{19}(2220)$	$(56, 4^+)^d$	320-550	198.0	66.0	264
$G_{19}(2250)$	$(70, 3^-)^q$	290-470	849	0	849
$G_{39}(2400)$	$(56, 3^-)^q$	330-450	514	0	514
$H_{3,11}(2420)$	$(56, 4^+)^q$	300-500	254.3	254.3	508.6
$I_{1,11}(2600)$	$(70, 5^-)^d$	500-800	145.6	145.6	291.2
$K_{1,13}(2700)$	$(56, 6^+)^d$	350±50	225.8	75.3	301.1
$I_{3,13}(2750)$	$(56, 5^-)^q$	350-500	721	0	721
$K_{3,15}(2950)$	$(56, 6^+)^q$	330-700	879	0	879

TABLE 6 : Λ , Σ widths (in MeV); $\Gamma(1)$ and $\Gamma(2)$ are the results of putting the s-quark outside and inside the diquark respectively.

State	SU(6) $\otimes$ O(3)	$\Gamma(\text{expt.})$	$\Gamma(1)$	$\Gamma(2)$	$\Gamma(\text{tot})$
$S_{01}(1405)$	$(70, 1^-)^d$	50 ± 2	69.1	49.8	56.2
$D_{03}(1520)$	$(70, 1^-)^d$	15.6 ± 1	19.8	22.0	21.3
$P_{01}(1600)$	$(56, 0^+)^d$	50-250	90.6	81.6	84.6
$S_{01}(1670)$	$(70, 1^-)^d$	25-50	81.6	63.0	69.0
$D_{03}(1690)$	$(70, 1^-)^d$	50-70	24.2	26.9	26.0
$S_{01}(1800)$	$(70, 1^-)^q$	200-400	178	109.5	132.4
$P_{01}^*(1810)$	$(70, 0^+)^d$	50-250	99.2	91.5	94.1
$F_{05}(1820)$	$(56, 2^+)^d$	70-90	31.5	33.8	33
$D_{05}(1830)$	$(70, 1^-)^q$	60-110	92.2	64.6	73.8
$P_{03}^*(1890)$	$(70, 0^+)^q$	60-200	103.4	76.1	85.2
$G_{07}(2100)$	$(70, 3^-)^d$	100-250	40.8	42.8	42.2
$F_{05}(2110)$	$(70, 2^+)^q$	150-250	193	115.8	141.5
$H_{09}(2350)$	$(56, 4^+)^d$	100-250	48.6	50.1	49.6
$P_{13}(1385)$	$(56, 0^+)^q$	36-39	42.1	37.1	38.8
$P_{11}(1660)$	$(56, 0^+)^d$	40-200	93.1	84.8	87.6
$D_{13}(1670)$	$(70, 1^-)^d$	40-80	23.6	26.3	25.4
$S_{11}(1750)$	$(70, 1^-)^q$	60-160	160.7	99.6	112
$D_{15}(1775)$	$(70, 1^-)^q$	105-135	83.5	59.3	67.4
$F_{15}(1915)$	$(56, 2^+)^d$	80-160	34.6	37.2	36.3
$F_{15}(1915)$	$(56, 2^+)^q$	80-160	189.8	88.4	122.2
$D_{13}(1940)$	$(70, 1^-)^q$	150-300	212	102.5	139
$F_{17}(2030)$	$(56, 2^+)^q$	150-200	119.8	77.7	91.7