PHYSICAL REVIEW D **74,** 094014 (2006)

Quarks, diquarks, and QCD mixing in the N^* resonance spectrum

Qiang Zhao[1,2,*] and Frank E. Close[3,†]

[1]*Institute of High Energy Physics, Chinese Academy of Sciences, Beijing, 100049, People's Republic of China*
[2]*Department of Physics, University of Surrey, Guildford, GU2 7XH, United Kingdom*
[3]*Rudolf Peierls Centre for Theoretical Physics, University of Oxford, Keble Rd., Oxford, OX1 3NP, United Kingdom*
(Received 10 March 2006; published 16 November 2006)

We identify a "Λ selection rule" for N^* resonances in the presence of QCD mixing effects. We quantify these mixing effects from existing data and predict amplitudes for exciting **20** representations in SU(6), which are forbidden in strict diquark models. By classifying Particle-Data-Group (PDG) states at $N = 2$, we show that $\gamma N \to K\Lambda$, $K^*\Lambda$, $K\Sigma$, $K^*\Sigma$, and $J/\psi \to \bar{p}N^*$ are ideal probes of baryon dynamics and for establishing whether strongly correlated diquarks survive for $L > 0$.

DOI: 10.1103/PhysRevD.74.094014 PACS numbers: 12.39.−x, 13.60.−r

I. INTRODUCTION

It is remarkable that 40 years after the quark model was first applied to the problem of baryon resonances [1] it is still not well established whether three constituent quarks are the minimal effective degrees of freedom or whether a quark-diquark dynamics, where a pair of quarks is "frozen" into a ground state, suffices. Indeed, the vast literature inspired by the apparent discovery of a metastable "pentaquark" baryon in 2003 [2] showed both how little the strong dynamics of quarks is understood and raised renewed speculation about the role and existence of highly correlated diquarks [3,4]. Furthermore, an acceptable description of baryon resonance spectroscopy has been proposed based on a quark-diquark picture [5].

The relative coordinate between the two quarks forming the diquark is constrained to be in the $l_\rho = 0$ state. (We use the symbols ρ, λ as in Ref [6] to denote the antisymmetric and symmetric two-body substates within a three-body wavefunction). In such a case, the familiar and established $[\mathrm{SU}(6), L^P]$ multiplets $[\mathbf{56}, L^P]$ and $[\mathbf{70}, L^P]$ occur, but it is impossible to form $[\mathrm{SU}(6), L^P]$ correlations, $[\mathbf{20}, L^P]$ [7]. Within a qqq dynamics, where both λ and ρ spatial oscillators can be excited, the spectrum is richer and such **20** states can occur.

Whereas $[\mathbf{56}, L^P]$ and $[\mathbf{70}, L^P]$ excitations are well established, the search for $[\mathbf{20}, L^P]$ has been largely ignored, primarily because they cannot be excited by mesons or photons from a **56** nucleon. This is because photons and $q\bar{q}$ beams transform under SU(6) as **35**, which with the SU(6) forbidden transition $\mathbf{56} \otimes \mathbf{35} \not\to \mathbf{20}$ causes them to decouple from nucleons in naive SU(6) [7,8].

The purpose of this paper is to re-examine the assumptions underlying resonance production in the quark model. This will lead us to a selection rule, that appears to have been overlooked in the literature, and also to identify circumstances where **20** states can be excited.

A standard and phenomenologically successful assumption common to a large number of papers in the quark model is that photon transitions are additive in the constituent quarks [9–11]. This assumption also underlies models of hadronic production and decay in the sense that when $q_1q_2q_3 \to [q_1q_2q_i] + [q_3\bar{q}_i]$, the quark pair q_1q_2 are effectively spectators and only q_3 is involved in driving the transition. Such approximations lead to well known selection rules, which have proved useful in classifying resonances [11]. We adopt this approximation as a first step and show that within it there is a further selection rule that appears to have been overlooked in the literature. We shall refer to this as the "Λ selection rule" and show how it may help classify N^* resonances.

The above "spectator-hypothesis" for transition amplitudes will be violated by the spin-dependent forces that act between pairs of quarks and break SU(6), such as those generating the N-Δ mass gap. For example, when the nucleon is in an electromagnetic field, gauge invariance and the presence of such two-body forces imply that there occur diagrams where the photon interacts with quark number 3, say, and the exchange force acts between quarks 3 and 1 or 2. Thus electromagnetic interactions can transfer momentum to both λ and ρ oscillators, which both spoils the spectator-hypothesis and opens the possibility of exciting **20**-plets.

As a specific and quantifiable example we shall assume these spin-dependent forces arise from gluon exchange in QCD. This has considerable quantitative support [12] and also has been shown to induce mixings, including **70**-plet configurations, into the nucleon wavefunction [13]. Taking into account that the N-Δ mass gap of 300 MeV is on the scale of Λ_{QCD}, the resulting mixing effects can be sizeable. Whereas $\mathbf{56} \otimes \mathbf{35} \not\to \mathbf{20}$, the coupling $\mathbf{70} \otimes \mathbf{35} \to \mathbf{20}$ is allowed. Consequently the SU(6) breaking that induces **70** correlations into the nucleon enables the excitation of **20**-plets by photons and mesons. Interestingly, we shall

*Electronic address: Qiang.Zhao@surrey.ac.uk
†Electronic address: F.Close@physics.ox.ac.uk

see that the Λ selection rule still manifests itself in $N^* \to K\Lambda$ transitions even though the SU(6) symmetry is broken. This phenomenon will be useful for clarifying the excitation of **20**-plets in the N^* spectrum.

In the next section we first present the "Λ selection rule". We then show how QCD-generated wavefunction mixing allows production of **20**-plets. We formulate these ideas in a QCD quark model [13] though the qualitative results should be more generally true. In the final section we quantify these effects and discuss their application to N^* classification.

II. THE MODEL

For reference, we specify our nomenclature. The standard SU(6) ⊗ O(3) wavefunction can be constructed from three fundamental representations of group S_3:

$$\mathrm{SU}(6)\!: \mathbf{6}\otimes\mathbf{6}\otimes\mathbf{6} = \mathbf{56}_s + \mathbf{70}_\rho + \mathbf{70}_\lambda + \mathbf{20}_a, \tag{1}$$

where the subscripts denote the corresponding S_3 basis for each representation, and the bold numbers denote the dimension of the corresponding representation. The spin-flavor wavefunctions can be expressed as $|\mathbf{N}_6, {}^{2S+1}\mathbf{N}_3\rangle$, where $\mathbf{N}_6$ ($= \mathbf{56}, \mathbf{70}$ or $\mathbf{20}$) and $\mathbf{N}_3$ ($= \mathbf{8}, \mathbf{10}$, or $\mathbf{1}$) denote the SU(6) and SU(3) representation and S stands for the total spin. The SU(6) ⊗ O(3) (symmetric) wavefunction is

$$|\mathrm{SU}(6)\otimes \mathrm{O}(3)\rangle = |\mathbf{N}_6, {}^{2S+1}\mathbf{N}_3, N, L, J\rangle, \tag{2}$$

where explicit expressions follow the convention of Isgur and Karl [14–16].

The basic rules follow from application of the Pauli exclusion principle to baryon wavefunctions together with an empirically well tested assumption that electroweak and strong decays are dominated by single quark transitions where the remaining two quarks, or diquark, are passive spectators [6]. In particular, selection rules for specific processes can resolve the underlying dynamics. For example, the Moorhouse selection rule [17] states that transition amplitudes for γp to all resonances of representation $[\mathbf{70}, {}^4 8]$, such as $D_{15}(1675)$, must be zero due to the vanishing transition matrix element for the charge operator.

Such correlations also lead to a "Λ selection rule", which appears to have been overlooked in the literature. It states that N^* in $[\mathbf{70}, {}^4 8]$ decouple from ΛK and ΛK^* channels. This follows because the $[ud]$ in the Λ has $S = 0$ and in the spectator approximation, the strangeness emissions in $N^* \to \Lambda K$ or ΛK^*, the spectator $[ud]$ in the N^* must also be in $S_{[ud]} = 0$, whereby such transitions for the N^* of $[\mathbf{70}, {}^4 8]$ with $S_{[ud]} = 1$ are forbidden.

Note that the Λ selection rule applies to both proton and neutron resonances of $[\mathbf{70}, {}^4 8]$, in contrast to the Moorhouse selection rule, which applies only to the proton. The nearest that we can find to this in the literature is that $\Lambda^*[\mathbf{70}, {}^4 8] \not\to \bar{K}N$ [18]. While the associated zero in $N^*[\mathbf{70}, {}^4 8] \to K\Lambda$ is implicit in work that has calculated the couplings of baryon resonances [19], the source and generality of the rule does not seem to have been noted [20].

III. APPLICATION TO N^* SPECTRUM

An immediate application of the rules is to the $D_{15}(1675)$, which is in $[\mathbf{70}, {}^4 8]$. According to the Moorhouse selection rule, the amplitudes for $\gamma p \to D_{15}$ should vanish. However, the experimental values are not zero, though they are small. Nonzero amplitudes arise from QCD mixings induced by single gluon exchange in the physical nucleon [13]. The effective interaction

$$H_{FB} = \frac{2\alpha_s}{3m_i m_j}\left[\frac{8\pi}{3}\mathbf{S}_i\cdot\mathbf{S}_j\delta^3(\mathbf{r}_{ij}) + \frac{1}{r_{ij}^3}\left(\frac{3(\mathbf{S}_i\cdot\mathbf{r}_{ij})(\mathbf{S}_j\cdot\mathbf{r}_{ij})}{r_{ij}^2} - \mathbf{S}_i\cdot\mathbf{S}_j\right)\right] \tag{3}$$

induces significant mixings between the ${}^2\mathbf{8}$ and ${}^4\mathbf{8}$ in the **56** and **70** [16] and the nucleon wavefunction becomes [13]

$$|N\rangle = 0.90|^2S_S\rangle - 0.34|^2S_{S'}\rangle - 0.27|^2S_M\rangle - 0.06|^4D_M\rangle, \tag{4}$$

where subscripts, S and M, refer to the spatial symmetry in the S and D-wave states for the nucleon internal wavefunction. Thus, the $O(\alpha_s)$ admixtures at $N = 2$ comprise a 34% in amplitude excited **56** and 27% **70** each with $L = 0$ and 6% **70** with $L = 2$. The **70** admixture quantitatively agrees with the most recent data [21] for the $\gamma p \to D_{15}$ amplitudes, neutron charge radius and $D_{05} \to \bar{K}N$ [13]. The results assume that mixing effects in the D_{15} are negligible relative to those for the nucleon [13]: this is because there is no $[\mathbf{70}, {}^2 8; L^P = 1^-]$ state available for mixing with the D_{15}, and the nearest $J = 5/2$ state with negative parity is over 500 MeV more massive at $N = 3$. Within this $O(\alpha_s)$ analysis, such mixing is negligible: transitions from the large components of the nucleon to small in D_{15} have $\Delta N = 3$. The leading $O(\alpha_s)$ amplitude for $\gamma p \to D_{15}$ is dominantly driven by the small components in the nucleon and the large component in the D_{15} [15] for which $\Delta N = 1$.

This violation of the Moorhouse selection rule supports the hypothesis of QCD mixing in the wavefunction of the N. The Λ selection rule also remains robust, in the context of the diquark model, as admixtures of $[ud]$ with spin one, which would violate it, are only expected at most to be 20% in amplitude [22], to be compared with 27% for the nucleon in Eq. (4). Therefore, we expect the Λ selection rule to be at least as good as the Moorhouse rule even at $O(\alpha_s)$. Thus decays such as $D_{15} \to K\Lambda$ will effectively still vanish relative to $K\Sigma$; for the $D_{15}(1675)$ the phase space inhibits a clean test but the ratio of branching ratios for the analogous state at $N = 2$, namely $F_{17}(1990) \to K\Lambda{:}K\Sigma$, may provide a measure of its validity. Secondly, for $\gamma n \to D_{15}$, where the Moorhouse selection rule does not apply, the amplitudes are significantly large and consistent with experiment [21]. However, due to the Λ selection rule, the $D^0_{15} \to K^0\Lambda$ which makes the search for the D_{15} signals in $\gamma N \to K\Lambda$ interesting. An upper limit of $B.R. < 1\%$ is set by the PDG [21] which in part may be due to the limited phase space; a measure of the ratio of branching ratios for $K\Lambda{:}K\Sigma$ would be useful. The $F_{17}(1990)$, which is the only F_{17} with $N = 2$, is an ideal candidate for such a test, which may be used in disentangling the assignments of the positive parity N^* at the $N = 2$ level.

The QCD admixture of $[\mathbf{70}, {}^2 8, 2, 0, 1/2]$ in the nucleon wavefunction enables the excitation of **20**-plets. There has been considerable discussion as to whether the attractive forces of QCD can cluster $[ud]$ in color $\bar{3}$ so tightly as to make an effective bosonic "diquark" with mass comparable to that of an isolated quark. Comparison of masses of $N^*(u[ud])$ and mesons $u\bar{d}$ with $L \geq 1$ support this hypothesis of a tight correlation, at least for excited states [5,23]. If the quark-diquark dynamics is absolute, then $SU(6) \otimes O(3)$ multiplets such as $[\mathbf{20}, 1^+]$ cannot occur. The spatial wavefunction for **20** involves both ρ and λ degrees of freedom; but for an unexcited diquark, the ρ oscillator is frozen. Therefore, experimental evidence for the excitations of the **20** plets can distinguish between these prescriptions.

Pauli symmetry requires an antisymmetric spatial wavefunction for a **20** e.g. for the lowest state $[\mathbf{20}, 1^+]$

$$\psi^a_{211}(\boldsymbol{\rho}, \boldsymbol{\lambda}) = \sqrt{2}(\rho_+\lambda_z - \lambda_+\rho_z)\frac{\alpha_h^5}{\pi^{3/2}} e^{-\alpha_h^2(\boldsymbol{\rho}^2+\boldsymbol{\lambda}^2)/2}, \quad (5)$$

where $\rho_\pm \equiv \mp(\rho_x \pm i\rho_y)/\sqrt{2}$ and $\lambda_\pm \equiv \mp(\lambda_x \pm i\lambda_y)/\sqrt{2}$, and hence both ρ and λ degrees of freedom have to be excited. Furthermore **20** decouples from $\mathbf{35} \otimes \mathbf{56}$ but is allowed to $\mathbf{35} \otimes \mathbf{70}$. A 27% **70** admixture in the nucleon has potential implications for resonance excitation that may be used to look for **20**-plets.

For transitions between representation $[\mathbf{70}, {}^2 8, 2, 0, 1/2]$ and $[\mathbf{20}, {}^2 8, 2, 1, J]$ the matrix element $\langle \psi^a_{210} | e^{ikr_{3z}} | \psi^\rho_{200} \rangle \equiv 0$. Nonzero photon transitions can occur by the orbital flip "electric" term. Since $L = 1$, the nucleon component $[\mathbf{70}, {}^2 8, 2, 0, 1/2]$ can be excited to $J = 1/2$ and $3/2$ corresponding to P_{11} and P_{13} in $[\mathbf{20}, {}^2 8, 2, 1, J]$. The helicity amplitudes (not including the mixing angle) are presented in Table I with the spatial integral

$$A = 6\sqrt{\frac{\pi}{k_0}}\mu_0 \frac{1}{g} \langle \psi^a_{211} | e^{ikr_{3z}} p_{3+} | \psi^\rho_{200} \rangle = -6\sqrt{\frac{\pi}{k_0}}\mu_0 \frac{1}{g} \times \frac{2k}{3\sqrt{3}} e^{-k^2/6\alpha_h^2}. \quad (6)$$

TABLE I. Helicity amplitudes for $[70, {}^2 8, 2, 0, 1/2] \to [20, {}^2 8, 2, 1, J]$ for P_{11} and P_{13}, and $[70, {}^2 8, 2, 0, 1/2] \to [70, {}^4 8, 1, 1, J]$ for D_{15}. A and B are the corresponding spatial integrals. The theory calculations with the mixing angle and experimental data [21] for $\gamma p \to D_{15}(1675)$ are also listed with unit $10^3 \times \text{GeV}^{1/2}$.

Final state	$A^p_{1/2}$	$A^p_{3/2}$
P_{11}	$-\frac{1}{3\sqrt{3}}A$	...
P_{13}	$\frac{1}{3\sqrt{6}}A$	$\frac{1}{3\sqrt{2}}A$
D_{15}	$\frac{1}{3}\sqrt{\frac{1}{15}}B$	$\frac{1}{3}\sqrt{\frac{2}{15}}B$
Theo.	14	20
Exp.	19 ± 8	15 ± 9

These amplitudes may be compared with those for $\gamma p \to D_{15}$ as listed in Table I, for which the spatial integral is

$$B = 6\sqrt{\frac{\pi}{k_0}}\mu_0 k \langle \psi^\rho_{110} | e^{ikr_{3z}} | \psi^\rho_{200} \rangle = 6\sqrt{\frac{\pi}{k_0}}\mu_0 k \times (-i)\frac{k}{3\alpha_h} e^{-k^2/6\alpha_h^2}. \quad (7)$$

Additional P_{11} and P_{13} from representation **20** automatically raise questions about the quark model assignments of the observed P_{11} and P_{13} states, among which $P_{11}(1440)$, $P_{11}(1710)$, and $P_{13}(1720)$ are well-established resonances, while signals for $P_{13}(1900)$ and $P_{11}(2100)$ are quite poor [21].

IV. POSITIVE PARITY N^* UP TO 2 GEV

At $N = 2$ in the quark model a quark-diquark spectrum allows $[\mathbf{56}, 0^+]$, $[\mathbf{56}, 2^+]$, and $[\mathbf{70}, 0^+]$ $[\mathbf{70}, 2^+]$. If all qqq degrees of freedom can be excited, correlations corresponding to $[\mathbf{20}, 1^+]$ are also possible.

This implies the following N^* and Δ states (the superscripts denoting the qqq net spin state as $2S + 1$):

$$\begin{aligned}
[\mathbf{56}, 0^+] &= P_{11}({}^2N); \qquad P_{33}({}^4\Delta) \\
[\mathbf{56}, 2^+] &= P_{13}, F_{15}({}^2N); \qquad P_{31}, P_{33}, F_{35}, F_{37}({}^4\Delta) \\
[\mathbf{70}, 0^+] &= P_{11}({}^2N); \qquad P_{13}({}^4N); \qquad P_{31}({}^2\Delta) \\
[\mathbf{70}, 2^+] &= P_{11}, P_{13}, F_{15}, F_{17}({}^4N); \qquad P_{13}, F_{15}({}^2N); \\
&\quad P_{33}, F_{35}({}^2\Delta).
\end{aligned} \quad (8)$$

Without the **20**-plets, a commonly accepted scheme is as follows [6,16].

The $P_{11}(1440)$ is assigned to $[\mathbf{56}, 0^+]$ at $N = 2$ with $P_{33}(1660)$ as its isospin 3/2 partner [24,25]. The photoproduction amplitudes off proton and neutron, $A^p_{1/2} = -0.065 \pm 0.004$ and $A^n_{1/2} = +0.040 \pm 0.010\ (\text{GeV})^{-1/2}$, are consistent with M1 transitions in ratio -3:2 as for the nucleons [9,10,21]. The mass splitting of the $N(1440) - \Delta(1660)$ is consistent with the hyperfine splitting.

For the $[\mathbf{56}, 2^+]$ the $F_{37}(1950)$ is uniquely assigned. The hyperfine mass splitting naturally associates the $F_{15}(1680)$ as a partner; this is further confirmed by its photoproduction amplitudes, which satisfy selection rules from both proton and neutron targets [9,10]. The $P_{13}(1720)$ is *prima facie* associated with the $F_{15}(1680)$ but the photoproduction amplitudes do not easily fit: the helicity 3/2 amplitude of the P_{13} is predicted to be one-half that of the F_{15}, which does not fit well with the data [21]. The overall message from the **56** is that the constituent quark degrees of freedom are manifested at $N = 2$ even though $N\pi$ couplings are strong as evidenced by $\Gamma_T \sim 350$ MeV for $P_{11}(1440)$. This sets the challenge of assessing the situation for the rest of the $N = 2$ levels.

Only three P_{11} states are expected at $N = 2$ if the **20**-plets cannot be excited. As to assigning these states: (i) The $P_{11}(1440)$ has already been assigned to **56**; (ii) the P_{11} of $[\mathbf{70}, {}^4 8; 2^+]$ is expected to be heavier [6,16] and the signals for $F_{17}(1990)$, $F_{15}(2000)$, $P_{13}(1900)$ and $P_{11}(2100)$ [21] lead to a plausible assignment of these four in $[\mathbf{70}, {}^4 8, 2, 2, J]$; (iii) this leaves $P_{11}(1710)$ a natural candidate for $[\mathbf{70}, {}^2 8, 2, 0, 1/2]$. Nonetheless, there are still states of $[\mathbf{70}, {}^2 8, 2, 2, J]$ missing, and another P_{13} and F_{15} are needed.

Within the approximations that excitation and decays are dominated by single quark transition, and that there is no mixing among the **56** and **70** basis states, the assignments of P_{13} and F_{15} can be determined. The Moorhouse and Λ selection rules give the following filters for these states: $P_{13}({}^4N)$ in $[\mathbf{70}, 0^+]$ decouples from $K\Lambda$ and γp but couples to γn. This contrasts with $P_{13}({}^2N)$ in either of $[\mathbf{70}, 2^+]$ or $[\mathbf{56}, 2^+]$ which couple to all three channels. The helicity 3/2 amplitude of F_{15} in either of the $[\mathbf{70}, 2^+]$ or $[\mathbf{56}, 2^+]$ couples to photons with amplitude that is twice the size of its P_{13} counterparts.

Table II shows the photoproduction amplitudes expected. The $P_{11}(1710)$ fits with $[\mathbf{70}, {}^2 8, 2, 0, 1/2]$ due to the large error bars, however neither of the $P_{13}(1710/1900)$ fit easily with being pure **56** or **70** states. When the QCD mixing effects are included the agreement improves, in that small couplings of $4N$ states to γp are predicted, in accord with data. However, the implication is the added complexity that an additional P_{11} and P_{13} correlation in $[\mathbf{20}, 1^+]$ is allowed. Most immediately this prevents associating the $P_{11}(1710)$ as $[\mathbf{70}, 0^+]$ simply on the grounds of elimination of alternative possibilities. Thus we now consider what are the theoretical signals and what does experiment currently say.

Qualitatively one anticipates $P_{13}({}^4N)$ having a small but nonzero coupling to γp, the γn being larger while the $K\Lambda$ decay is still forbidden. For the **20** states $P_{11,13}({}^2N)$ both γp and γn amplitudes will be small and of similar magnitude. However, mixing with their counterparts in **56** and **70** may be expected. In Table II, we list the helicity amplitudes for the $P_{11}(1710)$, $P_{13}(1720)$ and $P_{13}(1900)$ with all the possible quark model assignments and the mixing angles from Eq. (4). The amplitudes for the P_{11} and P_{13} of $[\mathbf{20}, {}^2 8, 2, 1, J]$ are the same order of magnitude as the Moorhouse-violating $\gamma p \to D_{15}(1675)$ in Table I. For the $P_{11}(1710)$ all three possible configurations have amplitudes compatible with experimental data. For $P_{13}(1720)$, assignment in either $[\mathbf{56}, {}^2 8; 2^+]$ or $[\mathbf{70}, {}^2 8; 2^+]$ significantly overestimates the data [27,28] for $A^p_{1/2}$ if it is a pure state. Table II shows that the presence of **20** cannot be ignored, should be included in searches for so-called "missing resonances", and that a possible mixture of the **20**-plets may lead to significant corrections to the results based on the conventional **56** and **70**.

As a benchmark for advancing understanding we propose the following scenario, as a challenge for experiment: can one eliminate the *extreme* possibility that $P_{11}(1710)$ and $P_{13}(1720)$ are consistent with being in **20** configurations? There are already qualitative indications that they are not simply **56** or **70**. Their hadronic decays differ noticeably from their sibling $P_{11}(1440)$: compared with the $P_{11}(1440)$ in **56** for which $\Gamma_T \sim 350$ MeV with a strong coupling to $N\pi$, their total widths are $\sim$100 MeV, with $N\pi$ forming only a small part of this.

Consider now the phenomenology were these states in **20**. As $\mathbf{20} \not\to \mathbf{56} \otimes \mathbf{35}$, whereas $\mathbf{20} \to \mathbf{70} \otimes \mathbf{35}$ is allowed, decay to $N\pi$ will be allowed only through the **70** admix-

TABLE II. Helicity amplitudes for the $P_{11}(1710)$ and $P_{13}(1720)$ with all the possible quark model assignments for them. The data are from PDG [21], and numbers have a unit of $10^3 \times \text{GeV}^{1/2}$. The numbers with "$*$" are from multichannel studies [26].

Reso.	Heli. amp.	$[\mathbf{56}, {}^2 8; 2^+]$	$[\mathbf{70}, {}^2 8; 0^+]$	$[\mathbf{70}, {}^2 8; 2^+]$	$[\mathbf{70}, {}^4 8; 0^+]$	$[\mathbf{70}, {}^4 8; 2^+]$	$[\mathbf{20}, {}^2 8; 1^+]$	Exp. data
$P_{11}(1710)$	$A^p_{1/2}$	$*$	32	$*$	$*$	-8	-15	$+9 \pm 22$
$P_{13}(1720)$	$A^p_{1/2}$	100	$*$	-71	17	7	-11	$+18 \pm 30$
	$A^p_{3/2}$	30	$*$	-21	29	12	-18	-19 ± 20
$P_{13}(1900)$	$A^p_{1/2}$	110	$*$	-78	14	9	-10	-17^*
	$A^p_{3/2}$	39	$*$	-28	24	16	-17	$+31^*$

tures in the nucleon. Using the wavefunction in Eq. (4), this implies that $N\pi$ widths will be suppressed by an order of magnitude relative to allowed widths such as from **56** or **70** initial states. Decays $\mathbf{20} \to N^*(\mathbf{70}) + \pi \to N\pi\pi$ or $N\eta\pi$ can occur in leading order when phase space allows. These results make the $P_{11}(1710)$ and $P_{13}(1720)$ extremely interesting as not only are their total widths significantly less than the $P_{11}(1440)$ but the dominant modes for $P_{11}(1710)$ and $P_{13}(1720)$ are to $N\pi\pi$, which allows a possible cascade decay of $\mathbf{20} \to N^*(\mathbf{70})\pi \to N\pi\pi$. In particular, $B.R._{P_{13}(1720)\to N\pi\pi} \sim 70\%$ could contain significant $N^*(\mathbf{70})\pi$; present partial wave analysis has not included this possibility [29]. For $P_{11}(\mathbf{20})$ the S-wave decay $P_{11} \to S_{11}\pi$ is expected to be a dominant two-body decay; hence we urge a search for $P_{11} \to N\eta\pi$. Decays to $\mathbf{56} + M$ (where M denotes a meson) will occur through mixing, as for photoproduction, and so $K\Lambda$ may be expected at a few percent in branching ratio. In general we advocate that partical wave analyses of $N\pi\pi$ allow for the possible presence of $N^*(\mathbf{70})\pi$.

In order to classify these states, we advocate that partial wave analysis quantifies the couplings to $N^*(\mathbf{70})\pi$ and also includes data on $K\Lambda$ and/or $K^*\Lambda$ relative to $K\Sigma$ as the Λ selection rule can then be brought to bear. The Λ selection rule is useful for classifying the P_{11} and P_{13} in either $[\mathbf{56}, {}^2 8]$ and $[\mathbf{70}, {}^2 8]$, or $[\mathbf{70}, {}^4 8]$ and $[\mathbf{20}, {}^2 8]$, by looking at their decays into $K\Lambda$ and/or $K^*\Lambda$. The Moorhouse selection rule can distinguish $[\mathbf{70}, {}^4 8]$ and $[\mathbf{20}, {}^2 8]$ since the $[\mathbf{70}, {}^4 8]$ will be suppressed in γp but sizeable in γn, while the $[\mathbf{20}, {}^2 8]$ will be suppressed in both. The $[\mathbf{70}, {}^4 8]$ decays to $K\Lambda$ will be suppressed relative to $K\Sigma$ for both charged and neutral N^*. $J/\psi \to \bar{p} + N^*$ is a further probe of N^* assignments, which accesses **56** in leading order and **70** via mixing while **20** is forbidden. Hence for example $J/\psi \to \bar{p} + (P_{11}{:}P_{13})$ probes the **56** and **70** content of these states. Combined with our selection rule this identifies $J/\psi \to \bar{p} + K\Lambda$ as a channel that selects the **56** content of the P_{11} and P_{13}.

In summary, with interest in N^* with masses above 2 GeV coming into focus at Jefferson Laboratory and accessible at BEPC with high statistic $J/\psi \to \bar{p} + N^*$, this new selection rule should be useful for classifying baryon resonances and interpreting $\gamma N \to K\Lambda$, $K^*\Lambda$, $K\Sigma$ and $K^*\Sigma$ [30,31]. A coherent study of these channels may provide evidence on the dynamics of diquark correlations and the presence of **20**-plets, which have hitherto been largely ignored.

ACKNOWLEDGMENTS

We are grateful to G. Karl for helpful comments. This work is supported, in part, by grants from the U.K. Engineering and Physical Sciences Research Council (Grant No. GR/S99433/01), and the Particle Physics and Astronomy Research Council, and the EU-TMR program "Eurodice", HPRN-CT-2002-00311, and the Institute of High Energy Physics, Chinese Academy of Sciences.

[1] R. H. Dalitz, Les Houches 1965, edited by C. DeWitt and M. Jacob (Gordon and Breach, New York, 1965), p. 251; Proceedings of XIII International Conf., Berkerly, 1966 (Univ. of California Press, Berkeley, CA, 1966), p. 215.
[2] T. Nakano *et al.*, Phys. Rev. Lett. **91**, 012002 (2003).
[3] R. Jaffe and F. Wilczek, Phys. Rev. Lett. **91**, 232003 (2003).
[4] M. Karliner and H. J. Lipkin, Phys. Lett. B **575**, 249 (2003).
[5] F. Wilczek, hep-ph/0409168.
[6] See S. Capstick and W. Roberts, Prog. Part. Nucl. Phys. **45**, S241 (2000), and references therein.
[7] M. Anselmino, E. Predazzi, S. Ekelin, S. Fredriksson, and D. B. Lichtenberg, Rev. Mod. Phys. **65**, 1199 (1993).
[8] D. Faiman and A. W. Hendry, Phys. Rev. **173**, 1720 (1968).
[9] L. A. Copley, G. Karl, and E. Obryk, Nucl. Phys. **B13**, 303 (1969).
[10] R. P. Feynman, M. Kislinger, and F. Ravndal, Phys. Rev. D **3**, 2706 (1971).
[11] F. E. Close, A. Donnachie, and G. Shaw, *Electromagnetic Interactions and Hadronic Structure* (Cambridge Univ. Press, Cambridge, England, 2006); F. E. Close, *Introduction to Quarks and Partons* (Academic Press, New York, 1978).
[12] A. De Rujula, H. Georgi, and S. L. Glashow, Phys. Rev. D **12**, 147 (1975).
[13] N. Isgur, G. Karl, and R. Koniuk, Phys. Rev. Lett. **41**, 1269 (1978); **45**, 1738(E) (1980).
[14] N. Isgur and G. Karl, Phys. Lett. B **72**, 109 (1977).
[15] N. Isgur and G. Karl, Phys. Rev. D **18**, 4187 (1978).
[16] N. Isgur and G. Karl, Phys. Rev. D **19**, 2653 (1979); **23**, 817(E) (1981).
[17] R. G. Moorhouse, Phys. Rev. Lett. **16**, 772 (1966).
[18] R. Cashmore, A. J. G. Hey, and P. J. Litchfield, Nucl. Phys. **B95**, 516 (1975).
[19] S. Capstick and W. Roberts, Phys. Rev. D **58**, 074011 (1998).
[20] W. P. Petersen and J. L. Rosner, Phys. Rev. D **6**, 820 (1972); D. Faiman, Phys. Rev. D **15**, 854 (1977).
[21] S. Eidelman *et al.* (Particle Data Group), Phys. Lett. B **592**, 1 (2004).
[22] N. Isgur and G. Karl, Phys. Rev. D **21**, 3175 (1980).
[23] F. E. Close, Proceedings of the 58th Scottish Summer School in Physics, St. Andrews, 2004 (unpublished).
[24] L. D. Roper, Phys. Rev. Lett. **12**, 340 (1964).

[25] R. H. Dalitz and R. G. Moorhouse, Phys. Lett. **14**, 159 (1965); **14**, 356(E) (1966).
[26] G. Penner and U. Mosel, Phys. Rev. C **66**, 055212 (2002).
[27] F. E. Close and Z.-P. Li, Phys. Rev. D **42**, 2194 (1990).
[28] Z.-P. Li and F. E. Close, Phys. Rev. D **42**, 2207 (1990).
[29] D. M. Manley and E. M. Saleski, Phys. Rev. D **45**, 4002 (1992).
[30] I. Hleiwaqi and K. Hicks, nucl-ex/0512039.
[31] L. Guo and D. P. Weygand, hep-ex/0601010.