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TESTING QCD WITH MILLI-TeV ELECTRONS

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TESTING QCD WITH MILLI-TOV ELECTRONS

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Analogies have been drawn between QCD where quarks build up hadrons, which in turn build up nuclei, and QED where electrons build up atoms, which in turn build up molecules. There has been a lot of talk about creating a quark theory of the nucleus in QCD by analogy with the way electrons build up molecules in QED. However, this seems to me to be a case of trying to run before you can walk, in that in QED the first step has been achieved, namely understanding atoms in terms of their electronic structure, but understanding molecules in QED is only understood in rather simple cases. In QCD the situation is even worse, in that we haven't even achieved the first step, namely in understanding how hadrons are built from quarks. We all believe that quarks are in the hadrons, and that at high momentum transfer we can apply perturbative QCD(pQCD). But to build the hadrons, we need pQCD and also non-perturbative QCD, which, at present, means building models to simulate the effects of the confinement. Insofar as any successful model of the nucleus will have it looking to a very good approximation as if it is made of nucleons, then I believe that first we have to understand the dynamics of quarks building the individual nucleons. Throughout this talk, I wanted to show how low-energy electron beams may help us deepen that understanding, and hopefully lead us towards the correct dynamical picture of the nucleon.

In order to set a framework for my talk, I will concentrate on the nonrelativistic constituent quark model (NRCQM). I choose this model because it has been around for 20 years or more, it is highly successful phenomenologically¹, has very little theoretical justification, and as a "true" description of nature, must be wrong. So progress in understanding could follow if first we make the model break down in experiments and then learn from the nature of its failure.

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I want to describe how photon interactions, such as provided at a milli-TeV electron accelerator, could provide important information in this regard. The photon is a clean probe, in that it couples to the spin and flavor of the constituent quarks and reveals the correlations among the flavors and spins inside the "dirty" target. The photon has spin one, so helicity amplitudes where the photon and target spin parallel (helicity three-half) or antiparallel (helicity one-half) give us a rich structure; and also since the photon has a mixture of isospin zero or one, both proton and neutron targets can give us rich information. By exciting nucleon resonances in various helicity states from proton or nucleon targets, we have enough amplitudes that ratios or linear combinations of them may wash the dirt away; by this I mean that some of the specific features of confinement which are common to all amplitudes may be factored out^{1,2,3}, and the remaining information which can be accessed by present theory may be probed. As a result, we can begin to learn some features of the detailed interactions among the quarks, and then as our confidence in modeling improves, we may begin to access the dynamics of the "dirt". By using the electron beam, we can create beams of virtual photons and tune their invariant mass; in this way, we can go to very virtual photons which probe the kinematic region where perturbative QCD applies, down to the real photon domain where the nonperturbative dynamics becomes important. In particular, I think we should cease to be so obsessed with the deep inelastic scattering data with very large Q^2 . We know that perturbative QCD is correct. There is now interest in seeing how Q^2 dependence beyond logarithms signals the onset of the nonperturbative regime.

There are already clues in the data that significant effects can be accessed by a milli-TeV electron machine. Furthermore, I shall show that current modeling is incomplete and/or inconsistent. So there are challenges both for theory and for experiment.

High-Energy Electron Production

There have been claims recently that a dramatic failure of the canonical quark model of nucleons has been manifested by deep inelastic scattering of polarized leptons from polarized protons. The claim is that quarks may carry none of the spin of the polarized proton⁴. If this is indeed true, then the quark model will be called into question and will at least have to be extended.

One interesting feature of these experiments is that the polarization asymmetry $A(x, Q^2)$

$$A \equiv \frac{\sigma_{1/2} - \sigma_{3/2}}{\sigma_{1/2} + \sigma_{3/2}}$$

(where $\sigma_{1/2, 1/2}$ are the photoproduction cross sections for transversely polarized photon and the polarized target to have net $\Delta_T = 1/2, 1/2$ appears to maximize in the kinematic limit of Bjorken $x \rightarrow 1$, i.e. $A^p(x \rightarrow 1) \rightarrow 1$, at least for a proton target (the only target so far studied). In modern QCD we believe that the asymmetry tending to unity as $x \rightarrow 1$ is due to chromomagnetic effects¹⁷ in the wave function (one gluon exchange). If the (yet to be measured) neutron asymmetry also maximizes in this limit, then we will have important confirmation of these ideas. These chromomagnetic effects are also manifested in low-energy experiments. Chromomagnetic forces split the $\Delta(P_{33}(1232))$ and nucleon masses — the well-known hyperfine splitting of energy levels. This already shows us that QCD can make clean-cut predictions for low-energy experiments. Less well known, perhaps, is that QCD perturbations induce mixing in the quark wave functions which may be accessed most sharply by electromagnetic probes. Furthermore, if the $x \rightarrow 1$ behavior of inclusive high-energy data relates to the low-energy exclusive behavior, $Q^2 \rightarrow \infty$, then we may look for some characteristic differences¹⁶ in the Q^2 dependence of Δ/N^* form factors, G_M^p/G_M^n — the elastic magnetic form factors, and other specific final states such as $ep \rightarrow e\Lambda K/e\Xi K$.

Spin-dependent sum rules imply that nontrivial Q^2 dependence must occur, probably in the low-energy region accessible to a milli-TeV machine. To illustrate this, consider the Drell-Hearn-Gerasimov (DHG) sum rule for real photons⁵

$$\frac{-2\pi^2\alpha}{M^2} \kappa^2 = \int \frac{d\nu}{\nu} (\sigma_{1/2} - \sigma_{3/2}). \quad (1)$$

(κ is the anomalous magnetic moment of the target mass M) and the Bjorken⁶ sum rule (for $Q^2 \rightarrow \infty$)

$$\frac{1}{6} \left| \frac{g_A}{g_V} \right| = \int_0^1 dx (g_1^p - g_1^n) \quad (2)$$

We can write these in a similar form if we define

$$G(\nu) \equiv \frac{m^2(\sigma_{1/2} - \sigma_{3/2})}{8\pi\alpha^2} \quad (3)$$

Thus the DHG sum rule becomes for protons

$$-\frac{1}{4} \kappa_p^2 = \int \frac{d\nu}{\nu} G^p(\nu, Q^2 = 0) \quad (4)$$

and for the proton-neutron difference

$$\int \frac{d\nu}{\nu} (G^p - G^n)(Q^2 = 0) = -\frac{1}{4} [\kappa_p^2 - \kappa_n^2] = 0.112 \quad (5)$$

The Bjorken sum rule (Eq. 2) becomes

$$\int_{Q^2/2m}^{\infty} \frac{d\nu}{\nu} (G^p - G^n)(\nu, Q^2) = \frac{1}{3} \frac{M^2}{Q^2} \frac{dA}{dQ^2} = \frac{1.37}{Q^2} \quad (6)$$

The EMC data for protons $\int_0^1 dx g_1^p(x, Q^2) = 0.114 \pm 0.029$ would suggest that

$$\int_{Q^2/2m}^{\infty} \frac{d\nu}{\nu} G^p(Q^2) = \frac{0.2}{Q^2} \quad (7)$$

Thus, we can compare the proton-neutron difference in Fig. 1 as a function of Q^2 . That the Q^2 dependence would be most interesting was noted long ago by Gilman Karliner and myself qualitatively^{2,7}, and the above quantification follows that of Ioffe and collaborators⁸.

Some comments about the significance of these figures is called for. Equation 2 applies in the scaling region. We know from unpolarized scattering that scaling is still rather good modulo logarithms at $Q^2 = 5 \text{ GeV}^2$. The logarithmic breaking of scaling only changes the curves in Fig. 1 very slightly and cannot make them approach the value for real photons. Indeed the curves in Fig. 1 have already overshoot the real photon value when Q^2 is as low as 4 GeV^2 . Thus, nonperturbative effects must enter at low Q^2 in order to make correspondence with the real photon value. How is this achieved? As yet, we do not

Fig. 1. Q^2 dependence of the Bjorken integral (Eqs. 6 and 7) for (a) p-n difference and (b) proton target assuming the validity of EMC data (Ref. 4). The values at $Q^2 = 0$ according to the DHG sum rule are indicated.

know. All of the above remarks have concentrated on the total invariant difference, but for proton targets alone, things appear to be even more interesting. As noted long ago⁷, there is an interesting change of sign for the polarization asymmetry A^p . For real photons, this asymmetry will be negative over some range of energy, whereas for highly virtual photons, the asymmetry is positive over a considerable range of x as shown by the recent data from EMC, confirming the old SLAC data⁹. While one might dispute errors in the EMC integral¹⁰, a very dramatic behavior would have to ensue as $x \rightarrow 0$ if the integral would be negative! Indeed, if that was the resolution of the above "paradox", it would, by itself, be most interesting! If the integral is of the order of 0.1 when Q^2 is greater than 4 GeV², then Fig. 1 shows that something significant must occur as one proceeds into the nonperturbative regime of small Q^2 physics. This is not just a slight smoothing effect; here nonperturbative effects seem to have to change coalitions radically! So with hints that something "has to give", combined with the fact that the integrals are weighted towards small energies which hints that the resonance region may be relevant in the sum rule, let's see what we know so far about these spin-dependent effects and then delineate the questions that arise for a milli-TeV electronic machine.

Photoproduction of N^* : Empirical

Above pion production threshold are three prominent resonance bumps. The first resonance is the $P_{33}(1232)$; the second consists of $S_{11}(1550)$ and $D_{13}(1520)$, and the third is mainly $F_{15}(1690)$. Although these are the most noticeable states, a glance at the particle data tables shows that there are many other resonances within these bumps. The electromagnetic couplings of many of these are not well known. A better knowledge of some of these could have a significant impact on theory.

The spin dependence of the above prominent resonances is known, at least qualitatively. The P_{33} is dominantly photoproduced by M1 radiation and so $A = -1/2$. If M1 dominates for all Q^2 as in the NRQM, this value of A will be preserved.

The D_{13} and F_{15} are both photoproduced from protons dominantly in the helicity 3/2 mode, hence $A = -1$. So we see that the prominent states at $Q^2 = 0$ have the right sign of A to help saturate the DHG sum rule. (However, the quantitative test of this sum rule, especially the contributions above the resonance region, is not yet done.)

One of the successes of the NRQM was its prediction², confirmed by data¹¹, that the polarization asymmetry would change sign rapidly with Q^2 . Thus, these resonances seem to do what is required to help resolve the conundrum mentioned in the previous section. However, there are interesting and significant

questions outstanding. Is the Q^2 dependence of D_{13} and F_{13} the same, or is A a function of Q^2 where ω is the resonance mass, or some other behavior? If A is a function of Q^2 ($\omega^2 - x^2$), then at what value of x^2 does A change sign and how does this correlate with the possibility that A change sign as $x \rightarrow 0$ in deep inelastic data? If A changes sign at fixed $Q^2 = 0.5 \text{ GeV}^2$, say, does this correlate with a possible similar behavior in the high W inelastic region? Interesting structure seems likely for all W and Q^2 ; we need to map it out to the best possible accuracy in order to understand the nonperturbative dynamics that are at work.

These spin dependences, and especially their Q^2 dependences, are among the sharpest probes of the nonperturbative dynamics. This is my challenge for experiment. Now I will address a parallel challenge for theory.

A Challenge for Theory

The early calculations of electromagnetic transitions in the NRQM used an interaction operator familiar in nuclear and atomic physics²⁶⁻²⁸

$$H_I = \sum_{i=1}^3 H_i; H_i \equiv \frac{e_i \vec{p} \cdot \vec{A}(r_i)}{m_i} + \mu_i \vec{\sigma} \times \vec{q} \cdot \vec{A}(r_i) \quad (8)$$

where quark i at position $\vec{r}_i$ has mass, charge, and magnetic moment m_i, e_i, μ_i , and $\vec{q}$ is the three-momentum of the photon. I will, for convenience, refer to these as the electric and magnetic interactions.

Thus, H was the sum of single-quark interactions. The electric term flips the L_z of the quark in the (SHO) potential; on absorbing a positive helicity photon, it transforms as L_+ . Similarly, the magnetic term flips the spin, transforming as S_+ . Thus, Eq. (8) transforms as

$$H = L_+ + \bar{u} S_+. \quad (9)$$

The most general form of the electromagnetic interaction for the sum of single-quark operators transforms as^{1,12}

$$H = L_+ + \bar{u} S_+ + c S_z L_+ + d S_- L_+ L_+ \quad (10)$$

(The fourth term only contributes for transitions to $L > 2$.) QCD, completely solved, would be needed to predict the relative weights $\bar{u}, c, d$; however, we can test the dynamics of the electromagnetic interaction by forming combinations of helicity amplitudes which eliminate these unknowns and depend only on the Clebsch-Gordan coefficients associated with the various L_+, S_z .

I will show that data imply that $c \neq 0$, hence that $S_z L_+$ contributes and that the operator in Eq. (1) is consequently insufficient. If one makes an expansion

of the relativistic electromagnetic interaction to $O(1/m^2)$, one finds a spin-orbit interaction with the desired form, 13,14 ,

$$H_{(1)} = \frac{e}{m} \vec{p} \cdot \vec{A} + u \vec{\sigma} \cdot \vec{q} \cdot \vec{A} + \frac{1}{2m} \left(2u - \frac{e}{2m} \right) \vec{\sigma} \cdot \vec{p} \times \vec{E} \quad (11)$$

where $\vec{p} \equiv -i\vec{\nabla}$ and $\vec{E}$ is the electric field. However, one can no longer write $H = 2H_1$ at this order: to separate the internal excitation from the center of mass necessitates additional "non-additive" terms in H — these terms are associated with the Wigner rotation of the quark spins in the recoiling system. Also, the binding potential, V_V or V_S (vector or scalar), enters at this order¹⁴, and these effects have not been consistently taken into account in the literature — indeed, there are even inconsistencies in the literature as to the nature of the correct H ¹⁵.

If we are to confront the data meaningfully, these issues must be better understood. I will set this as a challenge for theory and survey now the data and experimental possibilities.

Effects Arising from One-Gluon-Exchange

The large Q^2 inclusive data are summarized by the structure functions $F_{1,2}$ (unpolarized) and a polarization asymmetry

$$A \equiv \frac{\sigma_{1/2} - \sigma_{3/2}}{\sigma_{1/2} + \sigma_{3/2}} = \frac{g_1}{F_1}.$$

($\sigma_{1/2,3/2}$ denoting the photoproduction cross section where γ and N have net $J_z = 1/2, 3/2$). The structure function g_1 enters in the Bjorken sum rule.

In the naive $SU(6)$ quark parton model $F^N/F^P = 2/3$, $A^N = 0$, $A^P = 5/9$. In the CGK paper (Ref. 2) we showed that the set of coherent ("resonance") quantum numbers that correspond to these numbers are that all 56 and 70 states in $SU(6)$ contribute with equal strengths up to overall $SU(6)$ Clebsch-Gordan coefficients. This is useful to know because it enables one to see how the total cross section ratios and polarizations are affected if subsets of states have different Q^2 dependent dynamics. The following examples illustrate this.

There is the well-known phenomenon $F^N/F^P(x+1) \rightarrow 1/4$ which arises from perturbation of the internal spin structure of the nucleon^{1,16} such as from single gluon exchange¹⁷. This leads to the prediction that

$$A^N = A^P + 1; (x+1),$$

and there are hints that this may be realized by the polarized proton data. The

essential dynamics is that, as $x \rightarrow 1$ (or $Q^2 \rightarrow \infty$ for exclusives), u^* dominates (i.e., a flavor polarized parallel to the parent proton, or a flavor polarized parallel to a parent neutron).

These would imply that excitation of Δ^* (in particular, $\Delta(1232)$) should die out faster with Q^2 than does N^* . Moreover, all resonances where the quarks' spins are coupled pairwise to $S_D = 1$ will die out fast. These have consequences for the nucleon and for the $\gamma N \rightarrow \Delta$ transition.

$$\text{As } Q^2 \rightarrow \infty \begin{cases} G_M^N/G_M^P \rightarrow e_d/e_u \rightarrow -1/2 \\ E2(\gamma N \rightarrow \Delta(1232)) \neq 0 \end{cases}$$

A question is how high in Q^2 is required? Carlson has argued¹⁸ that the necessary Q^2 may be quite small — a few GeV^2 .

Finally, exclusive processes such as $\sigma(\gamma N \rightarrow K\pi)/\sigma(\gamma N \rightarrow K\Lambda)$ should die out with Q^2 as the ud pair in Σ have $S_D = 1$ in contrast to the $S_D = 0$ configuration in the Λ .

$P_{11}(1470)$: The Roper Resonance

In the NRQM this state is an $N=2$ oscillator excitation, $(56,0^+)^*$ dominantly, which is sandwiched between the $N=0$, $\Delta(1232)$ and the $N=1$ $D_{13}(1518)$ and $S_{11}(1550)$. It is interesting in view of its low mass — raising doubts about the NRQM interpretation (the possibility suggested in one analysis that it is, in fact, two states) and also interesting in that it has the P_{11} quantum numbers expected for the lightest hybrid (gluonic) baryon¹⁹. Thus, good data on this state's electromagnetic couplings are called for²⁰.

If it is an $N=2$ state, then its electromagnetic excitation form factor will be enhanced relative to the prominent $D_{13}(1518)$ due to the $N=2$ versus $N=1$ threshold factors. Consequently, the P_{11} may dominate²¹ over the D_{13} when $Q^2 > 1$ or 2 GeV^2 , even though it is suppressed when $Q^2 = 0$. This could be an example where the $Q^2 \neq 0$ facility inherent in electroproduction is an important lever. This relative "turning on" of the P_{11} may help a clean separation from the nearby $S_{11}(1550)$.

It is important to confirm whether the ratio of amplitude for $\gamma N \rightarrow P_{11}(1470)$ satisfy ($n_{1/2}$ refers to the helicity $1/2$ amplitude of neutron)

$$n_{1/2}/p_{1/2} = -2/3$$

as would be the case for magnetic dipole excitation of a $(56,0^+)^*$ state (analogous to the ratio of magnetic moments of the nucleon). Present data are consistent with this, and if this result can be tightened statistically, then it

would be the best confirmation of a clear $(56,0^+)^*$ constituency.

Furthermore, it would eliminate the possibility of it being the lightest hybrid baryon. There is a selection rule that states²

$$\gamma p \leftrightarrow (p\gamma); \quad \gamma n \leftrightarrow (n\gamma)$$

for the lightest P_{11} and P_{13} $I=1/2$ hybrids. Note that if the Roper is split²⁸, one state being hybrid and the other a $(56,0^+)^*$ NRQM state, then comparison of γN and πN data will be interesting. On proton targets two states would contribute to πN but only one to γN . The reversed process $\pi^- p \rightarrow N^* + \gamma n$ could access the neutron coupling of both NRQM and hybrid states.

Note also the important question of the existence of $P_{33}(1600)$. In a $(56,0^+)^*$ scenario the P_{11} will be partnered by P_{33} , as for the nucleon and $\Delta(1232)$. However, a hybrid P_{11} has no such partner. The status of the $P_{33}(1600)$ is far from clear. It should be searched for in γN data where it should be excited by M1 as is $\Delta(1232)$.

$D_{13}(1518)$ and $F_{15}(1690)$

These prominent members of the second and third resonances are well known, and I cite them as illustrative of the earlier remarks on the NRQM and also in their Q^2 dependence with relation to the polarization asymmetries.

In Table 1 I illustrate the range of data for the $Q^2 = 0$ helicity amplitudes and some model calculations. The fine details are not the main issue here. Note the global features that $p_{1/2}, n_{1/2} \ll p_{3/2}, n_{3/2}$, and that the NRQM has $p_{3/2} = -n_{3/2}$ and that QCD wavefunction mixings, as in Ref. 22, change this equality by only 5%. Empirically, however, $-n_{3/2} \approx 2/3 p_{3/2}$.

I have discussed this in detail in my textbook, so I will not repeat the details here. Suffice it to say that any model which restricts H to $L_+ + \bar{u}S_+$ (Eq.

Table 1. Comparison of data and quark models (restricted to the interaction form of Eq. 8) for $D_{13}(1520)$. Note the large $n_{3/2}, p_{3/2}$ amplitudes and that mixing (Ref. 23) only slightly disturbs the equality in their magnitudes.

	<u>Range of Data</u>	<u>FKR/CKO (Ref. 28)</u>	<u>KI (Ref. 23)</u>
$p_{1/2}$	-22 ± 10	$= 0$	-23
$p_{3/2}$	$+167$ to $+195$	$+110$	$+128$
$n_{1/2}$	-65 ± 13	-30	-45
$n_{3/2}$	-120 - 144	-110	-122

9) will, by virtue of the Clebsch-Gordan coefficients, be forced to have

$$n_{1/2} = -p_{1/2} \quad (12)$$

if the D_{13} is a pure ${}^2P_{3/2}$ state. One can also form the linear combination, elimination of the non-perturbative unknown $\bar{u}$, which reads

$$n_{1/2} + \frac{1}{3} p_{1/2} = -\frac{2}{3\sqrt{3}} p_{3/2}. \quad (13)$$

We are forced to have to go beyond the simple interaction. If we assume the most general form of single-quark excitation (Eq. 10) in the form $L_+ + \mu S_+ + RS_Z L_+$, then the above generalize to

$$n_{3/2} = -p_{3/2} \left(\frac{1+R/3}{1+R} \right) \quad (14)$$

$$n_{1/2} + \frac{1}{3} p_{1/2} = -\frac{2}{3\sqrt{3}} p_{3/2}/(1+R) \quad (15)$$

$Q^2 = 0$ data suggest that¹

$$\langle L_+ \rangle : \langle S_+ \rangle : \langle S_Z L_+ \rangle = 4:1:2 \quad (16)$$

and hence that the spin-orbit interaction cannot be ignored: at which point we have to face the breakdown of additivity (single-quark summation). Note that we can empirically evaluate this breakdown by eliminating R between Eqs. (14) and (15) and obtaining the single constraint

$$-\frac{1}{3} \left(p_{1/2} - \frac{1}{\sqrt{3}} p_{3/2} \right) = \left(n_{1/2} - \frac{1}{\sqrt{3}} n_{3/2} \right), \quad (17)$$

a failure of which would reveal the presence of non-additivity. Data are not good enough to test this definitively.

As we have seen, there are dramatic effects expected as Q^2 departs from photoproduction into the $Q^2 \neq 0$ domain. But, theorists note, no analysis of this Q^2 dependence has yet been made using a consistent electromagnetic operator.

The $F_{15}(1690)$ also gives rise to an interesting selection rule if only $L_+ + \mu S_+$ is employed²⁶. The F_{15} is $56L=2, S=1/2$ in the NRQM, and so the helicity $3/2$ excitation amplitudes require $L_z \neq 0$. As the nucleon has $L=0$, this amplitude requires L_z flip, hence the $e\vec{p} \cdot \vec{A}$ piece. But in flavor space this is proportional to

$$\langle 56 | e | 56 \rangle = +1 p; 0 n$$

and so $n_{3/2} = 0$.

This will become non-zero if there is Δ mixing, i.e., 2Ω components to the wavefunction in order to be only $\frac{1}{2}$. If Δ is non-zero, Δ $\psi_{\frac{1}{2},\frac{1}{2}}$ is present as this gives

$$\langle 56 | eS_2 | 56 \rangle_n = \mu_n \neq 0.$$

Precision data may help us evaluate to what extent single-quark excitation ("additivity") fails, and by the systematic (?) pattern of failure guide more clearly to the underlying dynamics.

For this state, too, there are non-trivial Q^2 dependences expected for the helicity amplitudes and spin asymmetry.

Hybrid Baryons

There are suggestions in the literature¹³ that the lightest "gluonic" or "hybrid" baryons (q^3g) are P_{11} and P_{13} states, possible below 2 GeV in mass. The $P_{11}(1470,1710)$, $P_{13}(1730)$ are the only real candidates at present. The photoproduction selection rule²⁰ seems to make $P_{11}(1470)$ more likely to be a $(56,0^+)$ conventional q^3 ; however, the $P_{11}(1710)$ photoproduction amplitudes are a problem in the NRQM and are similar to those anticipated for a hybrid state.

In Table 2 I show the amplitudes in the NRQM and data. We see that the absolute magnitudes are not good for $P_{11}(1470)$, but the ratio is consistent with $-2/3$, as befits M1 excitation of $(56,0^+)^*$. This is an example of how ratios or combinations of amplitudes reveal underlying dynamics.

The data on $P_{11}(1710)$ and P_{13} are less conclusive, but there are 3σ differences compared to the NRQM on protons. The $P_{11}(1710)$ may well have zero amplitude

Table 2. Data and models for P_{11}, P_{13} amplitudes. The first number in each box is datum as cited in the table in ref. 20; the second is the NRQM of ref. 23, and the third (in parentheses) is for hybrid baryon $P_{11}(1710)$ or $P_{13}(1730)$.

	$A_{3/2}^p$	$A_{1/2}^p$	$A_{3/2}^n$	$A_{1/2}^n$
$P_{11}(1470)$		-70 ± 10 -24		42 ± 20 16
$P_{11}(1710)$		-3 ± 15 -47 {0}		9 ± 30 -21 {36}
$P_{13}(1730)$	-34 ± 25 +46 {0}	60 ± 36 -133 {0}	-26 ± 80 -10 {35}	-6 ± 25 +57 {20}

as in the hybrid selection rule. We need improved data to establish this zero better. If it is a fundamental zero, as in a selection rule, then it will be true for all Q^2 , but if it is due to an "accidental" cancellation between various terms in the interaction, then it may well change with Q^2 . This is an example where Q^2 variation can give important information.

An explicit example of non-trivial Q^2 dependence is in the helicity structure, or polarization asymmetry, in the prominent D_{13} and F_{15} resonances discussed earlier.

Conclusion

In this brief survey, I hope to have shown that the field of resonance photo and electroproduction is rather fertile and that the progress in theory since the mid-1970's poses new challenges for experiment. The complete and consistent study of the electromagnetic transition operator (Eq. 8, et seq.) is in progress, and its consequences for the Q^2 dependence of resonance excitation, including QCD mixing effects, are being worked out. I hope that precision data will illuminate failings in the NRQM which, in turn, will direct us better into understanding the nature of confinement and non-perturbative QCD.

References

- 1) F.E. Close, Introduction to Quarks and Partons (Academic Press, New York, 1978). Many points in the present text are amplified there.
- 2) F.E. Close and F.J. Gilman, Phys. Letts. 38B (1972) 541; F.E. Close, F.J. Gilman, and I. Karliner, Phys. Rev. D6 (1972) 2533.
- 3) A.J.G. Hey and J. Weyers, Phys. Letts. 48B (1974) 69.
- 4) J. Ashman et al. (EMC collaboration), Phys. Letts. B206 (1988) 364.
- 5) S.D. Drell and A.C. Hearn, Phys. Rev. Letts. 16 (1966) 908; S.B. Gerasimov, J. Nucl. Phys. (USSR) 2 (1966) 598.
- 6) J.D. Bjorken, Phys. Rev. 148 (1966) 1467; D1 (1970) 1376.
- 7) F.E. Close, in: Procs. of IX Rencontre de Moriond, Vol. 2, ed. J. Tran Thanh Van (CNRS, France, 1974) p. 285.
- 8) B. Ioffe, V. Khoze, and L. Lipatov, Hard Processes, Vol. 1 (North-Holland, Amsterdam, 1984).
- 9) V.W. Hughes and J. Kuti, Ann. Rev. Nucl. Part. Sci. 33 (1983) 611.
- 10) F.E. Close and R.G. Roberts, Phys. Rev. Letts. 60 (1988) 1471.
- 11) See V. Burkert, Electroproduction of Nucleon Resonances at CEBAF Energies, CEBAF-R-86-011 (1986).
- 12) F.E. Close, H. Osborn, and A.M. Thomson, Nucl. Phys. B77 (1974) 241; F.J. Gilman and I. Karliner, Phys. Letts. 46B (1973) 426.

- 13) L.A. Copley, G. Karl, and E. Obryk, Nucl. Phys. B13 (1969) 303; H. Osborn, Phys. Rev. 176 (1968) 1514; ibid. (1968) 1523; S.J. Brodsky and J. Primack, Ann. Phys. 52 (1969) 315; F.E. Close and L.A. Copley, Nucl. Phys. B19 (1970) 477.
- 14) F.E. Close and H. Osborn, Phys. Letts. 34B (1971) 400; A. Le Youanc et al. Orsay report, 1988.
- 15) K. Ohta, Phys. Rev. Lett. 43 (1979) 1201.
- 16) F.E. Close, Phys. Letts. 43B (1973) 422.
- 17) F.E. Close and A.W. Thomas, Physics Letters (in print).
- 18) C. Carlson, CEBAF Workshop 1987; Phys. Rev. D34 (1986) 2704.
- 19) T. Barnes and F.E. Close, Phys. Letts. 123B (1983) 89; F. Golowich, E. Haqq, and G. Karl, Phys. Rev. D28 (1983) 160.
- 20) T. Barnes and F.E. Close, Phys. Letts. 128B (1983) 277.
- 21) T. Abdullah and F.E. Close, Phys. Rev. D5 (1972) 2332.
- 22) N. Isgur and G. Karl, Phys. Lett. 72B (1977) 109; 74B (1978) 353; Phys. Rev. D18 (1978) 4187; D19 (1979) 2653; D20 (1979) 1191; N. Isgur, G. Karl, and R. Koniuk, Phys. Rev. Lett. 41 (1978) 1269.
- 23) R. Koniuk and N. Isgur, Phys. Rev. D21 (1980) 1868.
- 24) R.P. Feynman, M. Kislinger, and F. Rerndal, Phys. Rev. D3 (1971) 2706.
- 25) Particle Data Group, Phys. Lett. 75B (1978) 1 as modified in table 7 of ref. 23.
- 26) D. Faiman and A. Hendry, Phys. Rev. 173 (1968) 1720; 180 (1969) 1572.
- 27) L.A. Copley, G. Karl, and E. Obryk, Nucl. Phys. B13 (1969) 303; Phys. Letts. 29B (1969) 117.
- 28) R. Arndt, J. Ford, and D. Roper, Phys. Rev. D32 (1985) 1085.