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THEORY OF LEPTON-HADRON INTERACTIONS*

C. H. Llewellyn Smith
Stanford Linear Accelerator Center, Stanford University, Stanford, California 94305

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In this talk I shall concentrate on three topics:

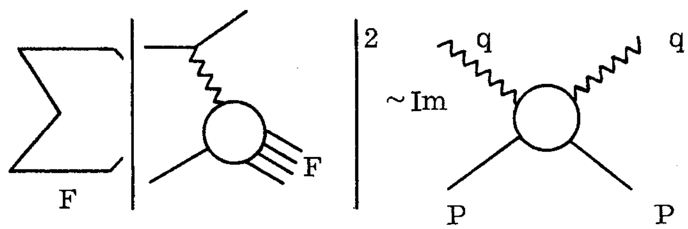
1. Partons, the light cone, etc.
2. Tests of models of hadrons in exclusive and inclusive electroproduction.
3. Unified, renormalizable(?) theories of weak and electromagnetic interactions.

1. Partons, The Light Cone, etc.

A. Relation Between the Parton Model and the Formal Light Cone Approach.

We shall consider:

$d\sigma^{eN \to e + \cdots} \sim$



$$\sim \int d^4y\, e^{iq\cdot y} < P \,|\, [J_\mu(y), J_\nu(0)] \,|\, P >$$

and the analogous neutrino reaction. We are interested in the (Bjorken) limit in which $\nu = q\cdot P$ becomes very large with $x = -q^2/2\nu$ fixed. For y^2 time like, the exponential factor oscillates extremely rapidly in this limit (unless $x \equiv 0$) except when y^2 is in the vicinity of the light cone (LC) $y^2 = 0$ (the commutator vanishes for y^2 spacelike). This oscillation wipes out the contribution from off the LC[1] (provided the matrix elements are reasonably behaved[2]). Therefore inelastic electron, and likewise neutrino, interactions measure the light cone commutator (LCC) of currents.

Taking spinless currents for simplicity, this LCC may be written[3]:

$$[J(y), J(0)]_{(y^2 \approx 0)} = C(y^2)\, J(y|0) + \text{less singular terms}$$

where C is a C number function which is singular on the LC and $J(y|0)$ a "bilocal operator" whose matrix elements are non-singular as $y^2 \to 0$. In free field theory J is given by products of fields, e.g. $\overline{\psi}(y)\psi(0)$. It can be shown formally that in many renormalizable interacting field theories the most singular part of the LCC is the same as in free field theory[4] (apart from a gauge factor $\exp(ig \int_0^y \phi_\mu(z)dz^\mu)$ if there is a vector interaction).

Formally has come to be an euphemism meaning "by methods which ignore the subtleties of field theory and are known to fail in perturbation theory". In the case of the LCC of currents the degree of singularity of C changes in perturbation theory (by a logarithmic factor)[5].

If the scaling observed in the SLAC-MIT experiments persists as $\nu \to \infty$ it would imply that C is the same as in free field theory (in contrast to perturbation theory in which C is such that scaling is broken by factors $\log(Q^2/M^2)$). Accordingly it is often assumed that perturbation theory is irrelevant and that the formal arguments are valid.

Making this very strong assumption, we can pick a field theory, set up the LCC, and make predictions (sum rules, relations between structure functions, etc.[6]). At this point we could abstract the Lorentz tensor and SU(3) properties of $J(y|0)$ and discard the field theory.[7] Alternatively, if we keep the field theory and make a Fock representation of the nucleon's state vector we see that, in the Bjorken limit, inelastic e and ν interactions are described by adding the contributions of the one body operators $\overline{\psi}(y)\psi(0)$; this is the parton model.[8]

Thus <u>in deep inelastic lepton scattering</u> the parton model is a particular realisation of an abstract light cone algebra. As long as we stick to relations which depend only on the Lorentz and internal symmetry properties of the model, the predictions of the LCC and the field theory (parton model) from which it is abstracted are obviously identical. The parton language can always be translated into the language of bilocal operators and may be employed even if we reject the idea that there is an underlying field theory.

Inelastic lepton scattering alone can therefore not establish the existence of a field theory (parton model). However, if there is such a field theory we can use it to investigate other processes which are not governed by the abstracted LCC. This has been done by Bjorken and Paschos, Drell, Levy and Yan, Landshoff and Polkinghorne, Brodsky, Close and Gunion, and others.[9]

An important application of parton models is to the process $PP \to \mu\bar{\mu} + \cdots$ which is described by

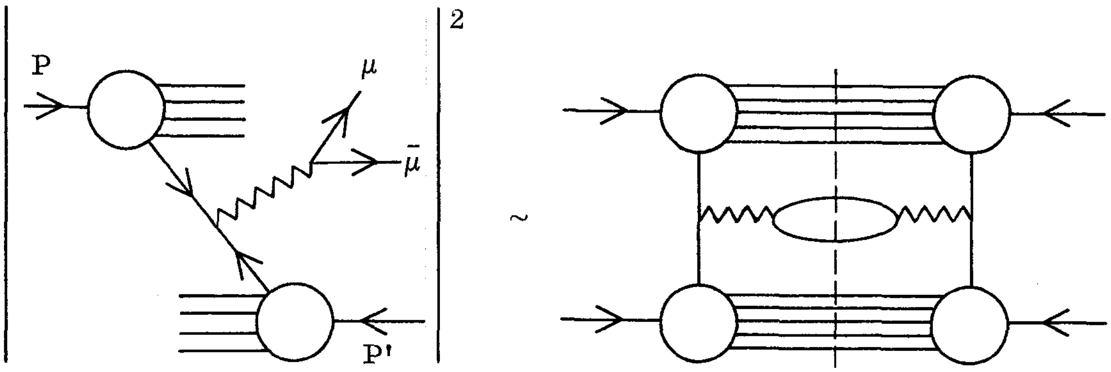
$$d\sigma \sim \int d^4y\, \Delta_+^R(y, Q^2) < PP' \text{in} \mid J_\mu(y) J_\nu(0) \mid PP' \text{in} >$$

where Q^2 is the invariant mass of the μ pair. We are interested in the limit $Q^2 \to \infty$. The cross section is LC dominated in the sense that the main contributions to the integral come from

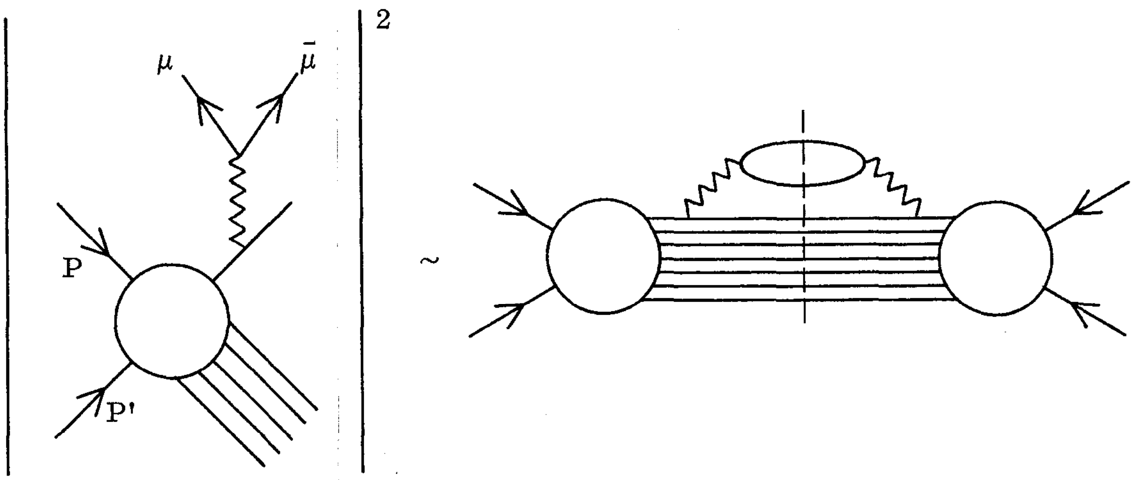
$$y^2 \lesssim \frac{1}{Q^2} \quad .$$

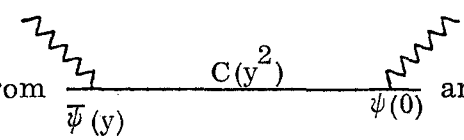
However, in contrast to inelastic lepton scattering, the leading LC singularity need not dominate. This is because the matrix element depends on another large parameter $S = (P+P')^2 > Q^2$ so that (e.g.) $S/Q^4 > 1/Q^2$, although a $1/Q^2$ term would be more singular by y^{-2} than an S/Q^4 term in coordinate space.

In fact, Drell and Yan found[10] that for finite Q^2/S, the dominant contributions in every order of their cut-off perturbation theory come from parton-antiparton annihilation:



and are not singular on the LC.[11] Nevertheless these contributions dominate those of the leading LC singularity, which correspond to Bremsstrahlung diagrams:



(the singularity comes from 

and is due to the parton propagating into the final state, just as in inelastic lepton scattering). Brandt and Preparata have contended[12] that the leading LC singularity dominates using Regge-like arguments. These arguments are presumably valid when $S >> Q^2$, in which regime the perturbation theory analysis is spoilt by the wee partons. Their results are therefore not necessarily incompatible with those of Drell and Yan.

In any case, the results obtained by Drell and Yan rely on the existence of a field theory and cannot obviously be attributed to some abstracted entity on general grounds. They predict that the differential cross section should satisfy the scaling law suggested by dimensional analysis

$$\frac{d\sigma}{dQ^2} = \frac{1}{Q^4} f\left(\frac{S}{Q^2}\right) .$$

The failure of this prediction would exclude parton models.

To be in a position to definitely establish the parton model we would like to calculate $f(S/Q^2)$. Unfortunately the value of f corresponding to the Drell-Yan diagram, which is given by folding parton into antiparton distributions (measured in eN and νN reactions), is presumably modified by the scattering of the spectator wees, although the basic interpretation and the scaling law should survive (actually Landshoff and Polkinghorne find[9] that the modification is small except when Q^2/S is small). However we could test the basic idea by comparing $PP \to \mu\bar{\mu} + \cdots$ with $\Pi P \to \mu^+\mu^- + \cdots$[13] or $\bar{P}P \to \mu^+\mu^- + \cdots$[14] for Q^2/S close to one. In this region the cross section is very sensitive to the (anti)parton distribution near $x = 1$; in most models the antiparton density in the proton is small near $x = 1$.[15] The ratio of $d\sigma(PP \to \mu\mu + \cdots)$ to $d\sigma$ for ΠP or $P\bar{P}$ should therefore decrease rapidly as Q^2/S approaches 1 if the model is correct (e.g., the cross sections for $P\bar{P}$ and PP differ by several orders of magnitude in this region in simple quark models[14]).

If the "parton sum rules" work in νN scattering, tests of these parton predictions for μ pair production could therefore then serve to establish whether this is really due to the existence of an underlying field theory.

B. Can Partons be Anything but Quarks?

It if is really true that the light cone commutator can be abstracted from free field theory, an immediate question is: what free field theory? Or more graphically (and, as discussed above, equivalently): what are partons?

First recall that in these models the electroproduction structure function νW_2 (e.g.) is given by:

$$\lim_{\substack{\nu \to \infty \\ x \text{ fixed}}} \frac{\nu W_2^{ep}}{M^2} = F_2^{ep}(x) = x \sum_i u_i(x) Q_i^2$$

where $u_i(x)$ is the probability of finding a parton of type i, which has charge Q_i, with a fraction x of the proton's momentum in the infinite momentum frame (more formally $u_i(x)$ is the matrix element of a piece of the bilocal; we employ the parton language below but all the properties used (e.g., $u_i \geq 0$) are true in both approaches).

We now consider several pieces of evidence which bear on the nature of partons:

a) σ_L/σ_T

It is by now well known that the small value of the longitudinal structure function F_L implies that the majority of the partons have spin $\frac{1}{2}$.[16] We assume below that $F_L \to 0$ in the Bjorken limit[17] and therefore all the charged partons have spin $\frac{1}{2}$.

b) Magnitude of $F_2^{ep, en}$

Energy momentum conservation gives:

$$1 = \int dx\, x \sum_i u_i(x) .$$

Hence if partons have integral charges:

$$\epsilon^{Q=0} \geq 1 - \int_0^1 F_2^{ep}\, dx = 0.83 \pm 0.02$$

where $\epsilon^{Q=0}$ is the fraction of the proton's momentum carried by neutral particles in the infinite momentum frame. This relation becomes an equality if $|Q| \leq 1$.

With a little more effort we can bound $\epsilon^{Q=Y=I=0}$ — the fraction of momentum carried by particles with $Q=Y=I=0$ — in many models. The results are given in the accompanying table.[19]

c) F_2^{ν}/F_2^{e}

Assuming chiral symmetry and CVC:

$$F_2^{\nu p+\nu n}(VV+AA) = 2F_2^{\nu p+\nu n}(VV) = 4\,(F_2^{ep+en})_{I=1} = R\,F_2^{ep+en} .$$

In models therefore:

$$R = \frac{F_2^{\nu p+\nu n}}{F_2^{ep+en}} = \frac{4(F_2^{ep+en})_{I=1}}{F_2^{ep+en}} = \frac{4\sum_i u_i(x)(I_z^2)_i}{\sum_i u_i(x)\left(I_z^2 + \frac{Y^2}{4}\right)_i} \leq \left(\frac{I_z^2}{I_z^2 + \frac{Y^2}{4}}\right)_{max.}$$

From this formula we immediately deduce the limits on R given in the table for various models.[20]

TABLE

Predictions and Restrictions in Various Models[18]

Model	$\epsilon_{Q=Y=I=0}$	$R \equiv \frac{F_2^{\nu p+\nu n}}{F_2^{ep+en}}$	$\frac{F_2^{en}}{F_2^{ep}}$	Remarks
Sakata Fermi-Yang	$=0.79\pm0.04$	$= 2$	anything	Special (only?) case of next category
Q integral $I \leq \frac{1}{2}$ $Q = Y/2 + I_z$	$\geq 0.79\pm0.04$	≤ 2	anything	
Σ Λ N Ξ	$\geq 0.79\pm0.04$	≤ 4 Chiral Symm.	anything	Need Σ's in nucleon to make $R > 2$
Q integral $Q = Y/2 + I_z$	$\geq 0.79\pm0.40$	≤ 4 Chiral Symm.	anything	Need $I \gtrsim Y$ partons in nucleon to make $R > 2$
Quarks G-M, **Z**	$\geq 0.52\pm0.38$	$\leq \frac{18}{5}$	$\geq \frac{1}{4}$	
Han-Nambu	$\geq 0.76\pm0.15$	$\leq \frac{10}{3}$	$\geq \frac{1}{2}$	*irrelevant if produced* hadrons SU(3)'singlets[22]
Experiment	$\frac{\int F_2^{\nu p+\nu n}\,dx}{\int F_2^{ep+en}\,dx} \geq 3.9\pm1.4$		As $x \to 1$, $\frac{F_2^{en}}{F_2^{ep}}$ appears to be between $\frac{1}{4}$ and $\frac{1}{2}$?	

(Note: The Han-Nambu model is the version used by Budny, Chang and Choudhury[23] who assumed that the nucleon is an SU(3)' singlet (thus ensuring that the usual quark model predictions obtain for the nucleon) but that states produced in electroproduction are not singlets, so that Lipkin's theorem[22] does not apply.)

d) F_2^{en}/F^{ep}

In some models the ratio F_2^{en}/F_2^{ep} is bounded, e.g. in the quark model[21].

$$4 \geq \frac{F_2^{ep}}{F_2^{en}} = \frac{\frac{4}{9}(u_p + u_{\bar{p}}) + \frac{1}{9}(u_n + u_{\bar{n}} + u_\lambda + u_{\bar{\lambda}})}{\frac{4}{9}(u_n + u_{\bar{n}}) + \frac{1}{9}(u_p + u_{\bar{p}} + u_\lambda + u_{\bar{\lambda}})} \geq \frac{1}{4}$$

where u_p, u_n ... are the distributions for the various species of quarks and we have used charge symmetry to equate u_p for proton targets to u_n for neutron targets, etc. Preliminary data presented at the Cornell conference suggested that the n/p ratio might fall below 1/4 but this now appears unlikely.[17] However it seems that it may fall below the lower limit of 1/2 allowed by the Han-Hambu model.

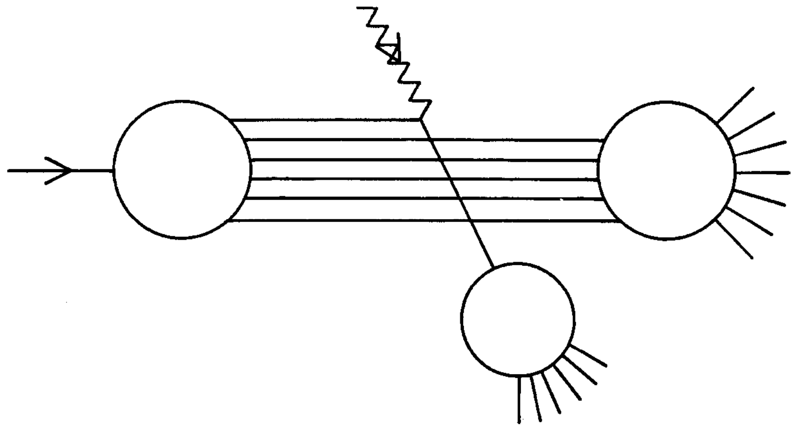
From the table we conclude that in models in which the partons have conventional quantum numbers the nucleon's momentum must mostly be carried by particles with $Q = Y = I = 0$ (and a fraction of the partons must have $I \gtrsim Y$). Such models would perhaps be a little contrived.

The quark model still awaits a definitive test (the F_3 sum rule will be crucial; it is very model dependent since it measures the average baryon number per interacting parton and is the easiest sum rule to test[6]). The quark model is at present compatible with all the inequalities[24] and, furthermore, simple explicit models work surprisingly well.[25]

C. Can Partons be Quarks?

Must Partons be Produced?

The parton model is based on diagrams with the structure:



This implies that if partons are quarks then the final state contains particles with non-integral baryon number. This separation of the produced particles into two distinct groups certainly obtains in perturbation theory (however, it has been argued that this could be changed by the "wee partons" which are not clearly associated with either group and are incorrectly treated in finite order[26]).

This result is also implied by the LCC if the bilocal can be factored into two operators with quark quantum numbers.[27] Inserting a complete set of states between these operators, and assuming that the asymptotic states are complete, the structure functions are found to vanish below the threshold for producing quarks.[22] Thus, unless quarks are produced, the applicability of the light cone algebra of the quark model seems to require that:

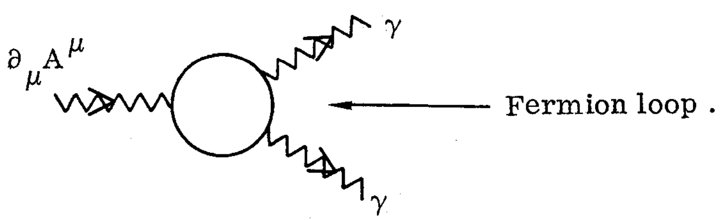
either: the bilocal does not factorise into two quark pieces (and the world is not described by the results of a formal treatment of quark field theories);

or: the asymptotic states are not complete, in which case quarks can "exist" but not "get out" of the nucleon (this idea has frequently been entertained and does not manifestly violate any sacred principles[28]).

Despite these results, it is possible to imagine that, at least in non-relativistic models, quarks at the bottom of a deep (but finite) potential well might behave much as free particles when subjected to a sudden impulse, although the sum rules might not converge to the expected values below threshold in this case. However, no respectable relativistic models of this type are known, nor are there soluble models in which the quarks do not "get out"; how the predictions for μ pair production would fare in these models is unclear.

The Bell-Jackiw, Adler Anomaly

Straightforward application of PCAC is known to forbid the decay $\pi^0 \to \gamma\gamma$ for $m_\pi = 0$.[29] However, the argument fails for the diagram:



This is the anomaly.[30,31,5] It turns out that in all orders of perturbation theory this is the only anomalous diagram.[5] Hence $\Gamma(\pi^0 \to 2\gamma)$ can be calculated exactly for $m_\pi = 0$. The value given by this diagram clearly depends sensitively on the charges of the elementary spin $\frac{1}{2}$ hadrons in the theory (the anomaly is independent of the masses). If the spin $\frac{1}{2}$ fields are protons and neutrons the resulting $\Gamma(\pi^0)$ agrees with experiment. In the quark model $\Gamma(\pi^0)$ is a factor of nine too small.

It is hard to know how seriously to take this result.[32] In the quark model the pion must be a bound state so that perturbative results are not necessarily relevant even when they are true to all orders.

Generally the anomaly will occur if fields have canonical dimensions[33] but its absolute strength cannot be calculated without recourse to perturbation theory. Similarly, canonical arguments imply that $\sigma(e^+e^- \rightarrow \gamma \rightarrow \text{hadrons}) \sim \frac{C}{S}$ but we cannot calculate C non-perturbatively (in perturbation theory[34] $C \sim \sum_i Q_i^2$). The problem in both cases is that we need a vacuum expectation value; since the conserved quantum numbers of the vacuum are all zero, there is no finite vacuum expectation value that is known on general grounds (contrast $\sigma^{\nu N}$ for which the known values of the nucleon's I, Y and B lead to sum rules which set an absolute scale). Although we do not know how to calculate the absolute values of $\Gamma(\pi^0)$ and $\sigma(e^+e^-)$ non-perturbatively, Crewther has shown[35] rather generally that these two quantities can be related in chirally symmetric models.

We could of course obtain the absolute values of $\Gamma(\pi^0)$ and $\sigma(e^+e^-)$ by simply postulating that the C number (vacuum expectation value) part of operator products on the light cone is given by free field theory. It should be noted that this postuate goes beyond the usual light cone assumptions (it seems to require — among other things — that the physical and bare vacua are the same) and we might be reluctant to abstract it since the only model in which it is known to be true (perturbation theory) undoubtedly requires that partons are produced in e^+e^- annihilation.[36]

The quark enthusiast who accepts this postulate is forced to conclude that there are three triplets of (para) quarks[37] in which case the predictions for $\Gamma(\pi^0)$ and $\sigma(e^+e^-)$ are increased by factors of nine and three respectively.

Even without this postulate, Crewther's results[35] and $\Gamma(\pi^0)$ suggest that $\sigma(e^+e^-)$ should be enhanced by a factor of three relative to the perturbative result in the quark model (although $m_\eta = 0$ is presumably a bad approximation, the value of $\Gamma(\eta \to 2\gamma)$ is a problem, as it is for all models based on SU(3) symmetry).

D. Conclusions

We should again emphasise that if scaling persists a fundamental question will be — why? The idea that scaling is due to the bound state nature of the nucleon is appealing and has recently been investigated by Drell and Lee[38] (whose work is closely related to earlier studies by West and by Landshoff and Polkinghorne). Unfortunately there is no soluble model to bear out this conjecture.

Existing data are easily explicable by the quark model. If scaling persists and the quark sum rules turn out to work, immediate questions will be:

1) Is there an underlying quark field theory? μ pair production might shed light on this question.

2) If scaling is associated with a quark field theory, why is scaling observed below the quark threshold?

3) How does $\sigma(e^+e^- \to \text{hadrons})$ behave? Is it twice $\sigma(e^+e^- \to \mu^+\mu^-)$ as suggested by the magnitude of $\Gamma(\pi^0)$? Does this imply the existence of three triplets of quarks?

2. Tests of Models of Hadrons in Exclusive and Inclusive Electroproduction

Electro and neutrino production provide a unique opportunity to test models by varying the mass of the effective projectile (photon or current). Here we shall discuss several examples, beginning with a case in which data already exist.

A. Resonance Electroproduction in the Quark Model

The non-relativistic quark model has successfully described photoproduction of low-lying baryon resonances.[39,40] Some of the successes depend only on the SU(6) × 0(3) structure of the model (and could be abstracted from it), e.g. because

$$H_{\gamma} = -\vec{\mu}\cdot\vec{B} - \frac{Q}{M}\vec{P}\cdot\vec{A}$$

transitions between states which are supposed to belong to the same 0(3) multiplet (e.g. $\gamma + N \rightarrow \Delta(1236)$), and are therefore induced by the spin raising operator, are M1 transitions. However, the (successful) prediction that the $D_{13}(1520)$ and $F_{15}(1590)$ resonances are predominantly photoproduced in the helicity 3/2 state depends on the details of the dynamics. Explicitly, in the non-relativistic model of Copley, Karl and Obryk[39], the helicity 1/2 amplitude with a proton target is given by:

$$A^{P}_{\frac{1}{2}} \sim |\vec{q}|^2 - \alpha^2/g \qquad (D_{13})$$

$$A^{P}_{\frac{1}{2}} \sim |\vec{q}|^2 - 2\alpha^2/g \qquad (F_{15})$$

where $\vec{q}$ is the photon's three momentum (arbitrarily?) evaluated in the isobars' rest frames and α and g are constants which can be fixed from other considerations. It happens that $|\vec{q}(F_{15})|^2 \approx 2|\vec{q}(D_{13})|^2$ and reasonable choices of α and g give small $\lambda = \frac{1}{2}$ amplitudes for both resonances in agreement with experiment. In any case, the model allows both $\lambda = \frac{1}{2}$ amplitudes to be small.

However, Close and Gilman have pointed out [41] that such an accidental cancellation should not persist in electroproduction since $|\vec{q}|^2$ increases monotonically with q^2. In fact the model of Feynman, Kislinger and Ravndal[42] predicts (e.g.) for the D_{13} that the ratio $|A_{3/2}|^2/|A_{1/2}|^2$ decreases from more than 10 at $q^2 = 0$ to less than one at $|q^2| = 0.3$ GeV^2 and $\sim 1/10$ at $|q^2| = 1$ GeV^2. Recent experimental data on π^0 electroproduction from Daresbury[43] and DESY[44] give no indication of any change in the helicity structure in the D_{13} and F_{15} regions out to $|q^2| = 0.6$ GeV^2 and 0.5 GeV^2 respectively.

We are forced to conclude that the success of this "dynamical prediction" at $q^2 = 0$ was accidental. Close and Gilman point out that since the transitions depend only on two (spin and orbital momentum raising) operators there are always two relations between the four amplitudes ($\lambda = 1/2$ and $3/2$ for proton and neutron targets) which can be used to test the more fundamental assumption of $SU(6) \times 0(3)$ symmetry.

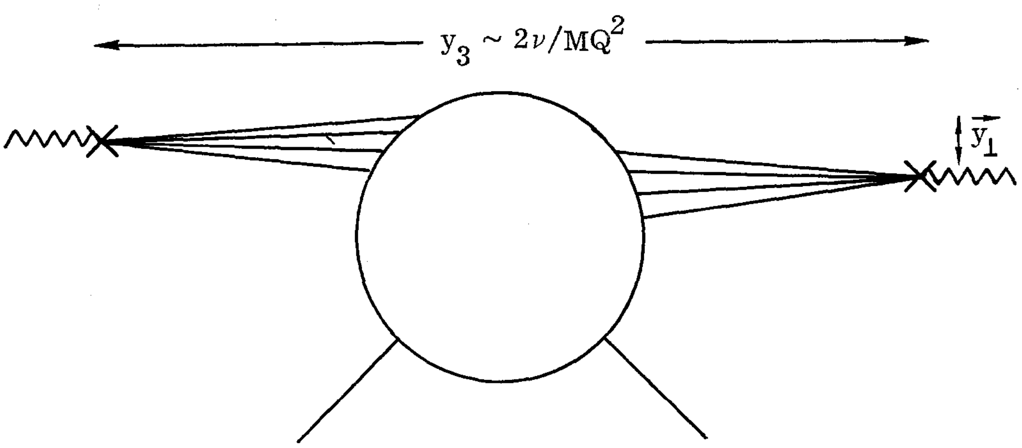
This provides an excellent example of the power of varying the photon's mass to test models. The moral is that theoretical descriptions of photoproduction should, when possible, be accompanied by predictions for electro (and neutrino) production, which may provide rather stringent tests.

B. Variable Photon Radius?

The monumental work of Cheng and Wu[45] shows that in Quantum Electrodynamics the photon's "radius" decreases when Q^2 increases. It has been suggested that this also occurs in the electroproduction of hadrons.[45,46] The argument for light cone dominance referred to in Section 1 [47] suggests that in forward Compton scattering of virtual photons at large ν:

$$\vec{y}_{\perp}^{\,2} \lesssim \frac{2M^2}{\nu^2\left(1-\sqrt{1+\frac{Q^2M^2}{\nu^2}}\right)} \quad \underset{(\nu \gg Q^2)}{\approx} \quad \frac{4}{Q^2}$$

In the diffractive region ($\nu \gg Q^2$) the longitudinal distance y_3 is much larger than the proton's radius and we can draw the following picture



It seems reasonable to infer that if $|\vec{y}_{\perp}|$ is small the hadronic "junk" into which the photon materialises must have a transverse position which is well defined to $\lesssim |\vec{y}_{\perp}|$ or, in other words, that

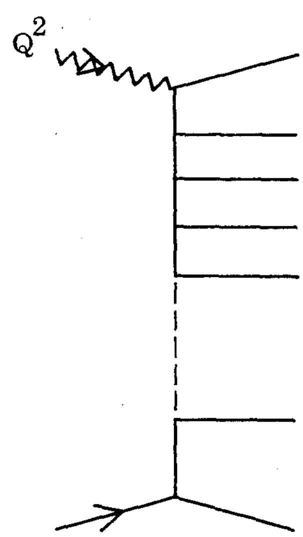
$$R_\gamma \lesssim \frac{1}{Q^2} .$$

In fact this argument suggests that at large ν, R_γ decreases with Q^2 both inside and outside the diffractive region. There is no firm agreement about this[48, 49]. In the absence of reliable hadronic models, we must await the verdict of experiment. The most crucial prediction is that the t distribution in $e + p \to e + p + \rho^0$ should broaden by up to a factor of 4 as Q^2 increases[50]; models suggest that significant changes should occur when Q^2 increases from 0 to 1 GeV^2.[48] Available data are contradictory.[17]

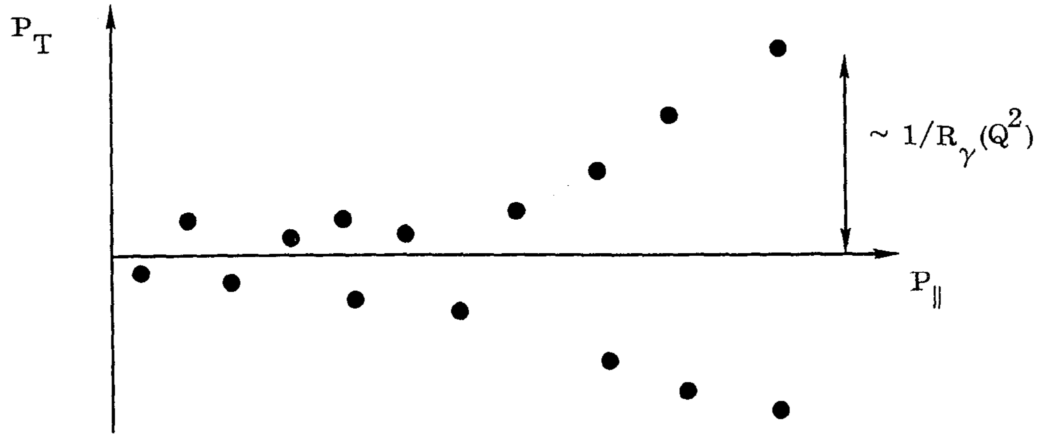
We are not here so much interested in this phenomenon in itself. Rather we intend to discuss tests of hadronic models which a variable R_γ would make possible, if it occurs. We emphasise that the following considerations are only relevant in regions in which $d\sigma_{ep \to ep\rho^0}/dt$ broadens as Q^2 increases, if this occurs at all.

Multiperipheral Model

Kogut has pointed out[48] (see also Nieh[49]) that a decreasing R_γ makes it possible to test the basic structure of the multiperipheral model:



Only the top few particles are senstive to Q^2. Recall that the fastest secondaries come from the top of the chain (the photon fragmentation region) and their transverse momentum distributions should broaden as R_γ decreases. However, $|\vec{P}_T|$ should be insensitive to R_γ for slower particles which emerge from further down the chain. Therefore a $\vec{P}_{\|}$, $\vec{P}_T$ plot for a given event should look like:



If R_γ decreases the absence of a broadening P_T distribution would cast severe doubt on the relevance of the multiperipheral model to electroproduction. The multiperipheral model (and more generally the concept of short range order) would also be in trouble if such a broadening persisted down to small values of $P_{\|}$ or, indeed, if the distribution of secondaries depended in any way on Q^2 except in the large $P_{\|}$ (photon fragmentation) region.

A Test of Dip Mechanisms

There are two schools of thought about dips in differential cross sections.

1. The "nonsense" school associates dips in $\Delta\lambda = 1$ amplitudes with nonsense points of the vector and tensor trajectories at $|t| \sim 0.5$ GeV^2.

2. The "geometrical" school argues that the imaginary parts of $\Delta\lambda = 1$ amplitudes are dominated by impact parameters $r \sim 1$ fm which are therefore described by "J_1"$(r\sqrt{-t})$, which has a zero at $|t| \sim 0.5$ for $r \sim 1$ fm.

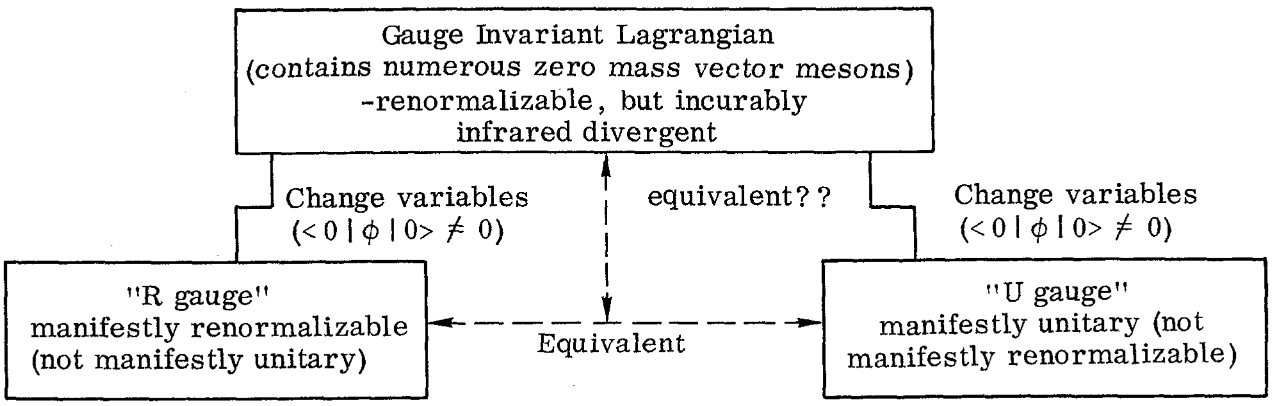
Harari has pointed out[51] that if R_γ decreases then the "geometrical school" would predict that in π^0 electroproduction the zero should move, e.g., if r^2 decreases by a factor of 2, the dip could perhaps move from $|t| \sim 0.5$ GeV^2 to $0.7 \lesssim |t| \lesssim 1$ GeV^2.[51,52] If the dip is due to a "nonsense" zero in the trajectory it would not move, since the properties of the exchanged trajectories are, of course, independent of the external particles. π^0 electroproduction may therefore provide a definite test of the origin of dips if it turns out that R_γ varies sufficiently.

3. Unified Renormalizable (?) Theories of Weak and Electromagnetic Interactions.[53]

"Gauge theories" of weak and electromagnetic interactions have recently excited a great deal of interest because

1. These theories are probably renormalizable.[54,55,56]
2. They unite weak and electromagnetic interactions in an elegant manner.[57,58,59,60]

In these theories there are massive vector mesons (W's) coupled to currents which involve massive fermions. Such theories contain highly divergent terms, because of the $k_\mu k_\nu / M^2$ piece of the propagators, and were generally believed to be unrenormalizable (e.g., the conventional "phenomenological" theory of weak interactions would require the introduction of an increasing number of arbitrary constants in each order to render all matrix elements finite). However, in the "gauge" theories the various masses and coupling constants are not independent and seemingly miraculous cancellations occur which make them renormalizable. These cancellations have their origin in the fact that the Lagrangian is derived from a primary ("stage 1") Lagrangian containing massless Yang Mills (gauge) fields; it is believed that the ultra-violet divergences can be renormalized in such theories.[61] Schematically the various stages in the theory are related thus:



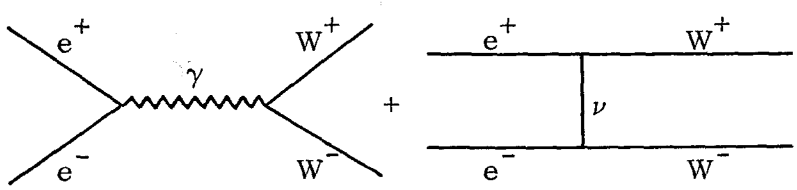
The gauge theory in stage 1 is renormalizable because the symmetries give rise to systematic cancellations of divergencies. However, the gauge invariance of the Lagrangian requires the vector mesons (gauge fields) to be massless. We need to retain the symmetrical Lagrangian but find a less symmetrical solution of the theory in which the vector mesons acquire mass, i.e., the symmetry must be spontaneously broken. Usually spontaneous symmetry breaking gives rise to (unobserved) massless mesons (Goldstone bosons). However in this case the conditions in which the Goldstone theorem holds do not obtain;[62,63,64] what happens is that the massless vector meson (with two degrees of freedom) and the incipient Goldstone boson (1 degree of freedom) conspire to form a massive vector meson (3 degrees of freedom)—this is known as the Higgs phenomenon.[62]

The theories are designed so that a U(1) gauge symmetry (of the type familiar in electro-dynamics) remains unbroken and there is therefore one residual massless neutral vector meson--the photon. The massive vector mesons mediate the weak interactions.

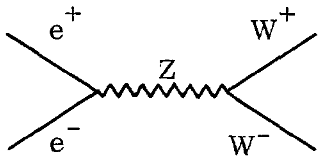
The spontaneous symmetry breaking is implemented in practice by introducing a multiplet (multiplets) of scalar mesons ϕ_i in stage 1 which are coupled to the vector mesons (and other particles) and have an invariant self interaction $V(\phi_i)$ whose minimum is not at $\phi_i = 0$ (so that $\langle 0 | \phi_i | 0 \rangle \neq 0$). A classical change of variables is performed so that the new fields vanish at the potential minimum (V has a minimum on a surface in ϕ_i space; the symmetry is broken by the choice of one particular point on the surface at which the new fields vanish). In the new variables it is obvious that some of the scalar mesons have become massive; the others are subsumed by the vector mesons which acquire mass.[65] The change of variables can be performed so that the resulting Lagrangian is manifestly renormalizable but seemingly non-unitary (the R gauge) or so that it is manifestly unitary but apparently unrenormalizable (the U gauge).

The recent resurgence of interest in these theories stems from 't Hooft's demonstration[54] that the S matrix is in fact unitary in the R gauge. B. W. Lee and Jean Zinn-Justin have recently shown[56] that the S matrix is actually the same in the R and U gauges (but it is probably not equivalent to the S matrix of the quantized stage 1 theory). [As suggested by Veltman, and demonstrated in detail by Gross and Jackiw,[66] and Bouchiat, Iliopoulos and Meyer[67] the renormalization is spoilt by triangle anomalies of the type discovered by Bell and Jackiw and Adler. However, only the triangle diagrams are anomalous,[68] and the contributions of different particles may cancell one another[66,67,69] if the theory is judiciously chosen.]

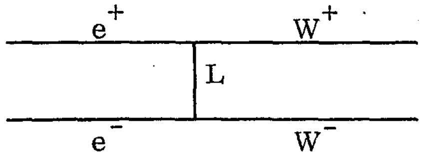
This is all very elegant but the question is--"does Almighty God in his omniscience make use of this rather clever scheme?"[70] If so, what is the relevant gauge symmetry? There is a physical argument which shows the possibilities rather clearly. Consider the process $e^+e^- \to W^+W^-$ which is described by the diagrams:



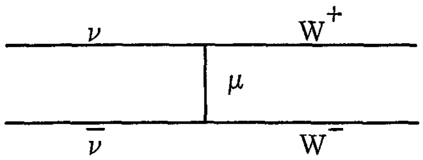
in the conventional theory. We can choose the $e^{\pm}$ helicities so that the second diagram does not contribute. The amplitude corresponding to the first diagram violates unitarity at high energies if the W's are longitudinally polarized. In order for the theory to make sense in perturbation theory at high energy, this badly behaved amplitude must be cancelled by other contributions in the same order.[71] There are two possibilities: either we must add a heavy vector meson(s) (Z) in the s channel:



or a heavy lepton(s) (L) in the t or u channels:



A similar argument applies to $\nu\bar{\nu} \rightarrow W^+W$. The amplitude corresponding to the diagram



is badly behaved at high energies. Again the situation may be saved by adding neutral vector mesons or heavy leptons (or both).

Thus we see that these theories require <u>either</u> the introduction of neutral Z's <u>or</u> heavy leptons (or both).[72] Weinberg has constructed a theory of the first type which has been widely discussed.[58,73,74,75] The symmetry is $SU(2)_L \times U(1)$. In stage 1 the fields which govern the leptonic interactions transform under $SU(2)_L$ as follows:

$$\begin{pmatrix} W_\mu^+ \\ W_\mu^o \\ W_\mu^- \end{pmatrix} \quad \text{- triplet.}$$

$$(1-\gamma_5)\begin{pmatrix}\nu_e\\ e\end{pmatrix},\quad (1-\gamma_5)\begin{pmatrix}\nu_\mu\\ \mu\end{pmatrix},\quad \begin{pmatrix}\phi^o\\ \phi^-\end{pmatrix},\quad \begin{pmatrix}\phi^+\\ \bar{\phi}^o\end{pmatrix} \qquad \text{- doublets}$$

B_μ, $(1+\gamma_5)e$, $(1+\gamma_5)\mu$ - singlets.

In stage 2, W^o_μ and B_μ mix yielding the photon and (together with ϕ) a heavy vector meson denoted Z_μ. $W^\pm_\mu$ combine with $\phi^\pm$ and acquire mass. There is one remaining massive neutral scalar meson. The parameters in the theory are two coupling constants g and g' and $\lambda = \langle 0 | \phi | 0 \rangle$. Because the theory is unified, the observed couplings and masses are dependent:

$$e = \frac{g g'}{\sqrt{g^2+g'^2}} \quad , \quad \frac{G_F}{\sqrt{2}} = \frac{g^2}{8 M_W{}^2} = \frac{1}{2\lambda^2} \; .$$

$$M_W = \frac{1}{2} g \lambda \geq 37 \text{ GeV.}, \qquad M_Z = \frac{1}{2}\sqrt{g^2+g'^2}\,\lambda \geq 75 \text{ GeV.}$$

The essential difference between this and the conventional phenomenological theory is the existence of the neutral Z which changes the predictions for processes such as $\nu_e + e \rightarrow \nu_e + e$.[75,76] The Z is a source of difficulties when hadrons are introduced; the data demand that it is essentially decoupled from the $\Delta S = 1$ hadronic current. The simplest scheme in which this occurs is the four quark model discussed by Glashow, Iliopoulous and Maiani[77] in which there are four quarks, whose left handed components from $SU(2)_L$ doublets:

$$(1-\gamma_5)\begin{pmatrix}p\\ n\cos\theta + \lambda\sin\theta\end{pmatrix}, \quad (1-\gamma_5)\begin{pmatrix}p'\\ n\cos\theta - \lambda\sin\theta\end{pmatrix},$$

the right hand components being singlets. The Z decouples from $\bar{n}\lambda$ because of a cancellation between the contributions of the two doublets. $\Delta S = 0$ neutral currents remain e.g., Weinberg estimates[74]

$$0.15 \leq \frac{\sigma(\nu + p \to \nu + p)}{\sigma(\nu + n \to \mu^- + p)} \lesssim 0.25$$

$$\frac{\sigma(\nu + p \to \nu + \pi^+ + n)}{\sigma(\nu + p \to \mu^- + \pi^+ + p)} \approx \frac{1}{9}.$$

(the limit 0.15 was obtained in an approximate calculation; treating the form factors carefully it is reduced to ~ 0.07 for reasonable choices of M_A).[78] Budny has shown[79] that, using a parton model;

$$\frac{\sigma^{\nu A \to \nu + \cdots} + \sigma^{\bar{\nu} A \to \bar{\nu} + \cdots}}{\sigma^{\nu A \to \mu^- + \cdots} + \sigma^{\bar{\nu} A \to \mu^+ + \cdots}} \gtrsim 25\% , \quad (A \sim p + n)$$

in this model. These results are on the verge of being excluded by experiment.

An interesting feature of this theory is that the triangle anomalies cancel internally if the quarks have the charge assignment: +1, +1, 0, 0 for p, p', n and λ.[53,67] In fact, there is a cancellation between the leptonic and the hadronic anomalies which requires the existence of the muon as well as the electron!

We now consider the model of Georgi and Glashow,[80] who choose the second alternative and introduce heavy leptons. The gauge symmetry is SU(2). There is a triplet of vector mesons, of which the neutral member becomes the photon in stage 2. The weak and electromagnetic charges of the known leptons do not generate an SU(2) algebra and additional leptons must therefore be introduced (the physical reason for this was discussed above). In stage 1 there are left and right handed triplets

$$\psi_L \sim \begin{pmatrix} X^+ \\ X^o \\ e^- \end{pmatrix}_L , \quad \psi_R \sim \begin{pmatrix} X^+ \\ X^o \cos\beta + \nu \sin\beta \\ e^- \end{pmatrix}_R ,$$

and corresponding triplets involving the muon (with heavy leptons denoted Y) and a triplet of scalar mesons. The remaining particles are singlets.

In stage 2 the theory has the following properties:

1. There are massive heavy leptons $X^{\pm}$, X^o, $Y^{\pm}$, Y^o and one massive scalar meson.
2. There are no neutral currents.
3. There are no triangle anomalies.
4. $M_W \lesssim 53$ GeV.
5. N_e and N_μ are conserved, M_ν remains zero in all orders but μ-e universality must be put in by hand.

Georgi and Glashow need 5 "quarks". A triplet:

$$\begin{pmatrix} p \\ n \\ \lambda \end{pmatrix}$$

with integral charges and two singlets: Q^- and Q^o. The strong interactions are U(5) invariant. In order that $M_{K_L^o} - M_{K_S^o}$ and $\Gamma(K_L^o \rightarrow \mu\bar{\mu})$ do not come out too large they need $M_W \lesssim 4$ GeV. Bjorken has pointed out that these second order effects can be suppressed by internal cancellations at the price of introducing 8 quarks.[81]

The demand that the contributions of the heavy leptons and the scalar meson(s) do not spoil the agreement between theory and experiment for g-2 for the muon puts severe constraints on these theories. It may already rule out the Glashow-Georgi theory, since it requires the "heavy" lepton to be very light and M_W to be close to its upper limit.[82]

To summarise the status of these theories:

1. They are probably renormalizable.

2. They involve either neutral currents or heavy leptons, which should be produced (e.g.) by sufficiently energetic neutrinos. In either case the theories can be tested experimentally.

3. The synthesis of weak and electromagnetic interactions is elegant. These theories may lead to further unifications of the particle spectrum; e.g., the absence of anomalies in the four quark model requires the existence of the muon as well as the electron--the interdependence of the lepton and the hadron spectra in this model is intriguing.

4. A major (aesthetic ?) drawback is that no one has yet succeeded in constructing a model which incorporates hadrons in a natural and elegant way compatible with the observed SU(3) structure of the spectrum.

4. Concluding Remarks

I have not had time to discuss the many fundamental questions which are still open experimentally (such as the existence of T violating or isotensor parts of the electromagnetic current, second class currents, etc.[83]). Nor have I discussed well established theoretical models. In fact most of this talk has been devoted to speculative ideas. It is perhaps necessary to emphasise this because, by dint of repetition, people tend to forget that many of the currently fashionable theoretical ideas are lacking an experimental basis. Perhaps scaling is broken by small $\log(Q^2/M^2)$ terms. Very likely the formal derivations of the light cone commutator are totally wrong. Indeed, it should be borne in mind that even the sacrosanct principles of SU(3) current algebra are still essentially untested (what if the Adler sum rule is wrong?). At present there is no experimental evidence for the existence of W's, let alone for gauge theories.

However, the theories I have discussed have one outstanding merit--they make predictions which can and will be tested. Perhaps by the time of the next conference in this series they will be established or destroyed.

Acknowledgements

I am grateful to my colleagues at SLAC for many useful discussions. In particular, I wish to thank J. Ellis and J. D. Bjorken for discussions related to sections 1C and 4 respectively and R. L. Jaffe who read the manuscript and made numerous helpful comments. I am also grateful to A. Pais for comments on the manuscript.

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1. This was pointed out by R. A. Brandt, Phys. Rev. Letters 23, 1260 (1969), B. L. Ioffe, Phys. Letters 30B, 123 (1969), L. S. Brown in Lectures in Theoretical Physics, ed. W. E. Britten, B. W. Downs and J. Downs (Interscience, New York, 1969) and doubtless several others.

2. R. L. Jaffe (SLAC-PUB-1023, 1972) has shown that singularities in the matrix elements on fixed mass hyperbolas would have to be more singular than those on the L. C. by at least $(X^2)^{-3/4}$ in order to give comparable contributions (which would not scale but be oscillating functions of Q^2 in the Bjorken limit).

3. R. A. Brandt and G. Preparata, Nucl. Phys. B27, 541 (1971).
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4. Perhaps the most elegant way to do this is to quantise the theory on a light like surface in which case the LCC of currents can be calculated in terms of the (canonical) LCC of fields just as the equal time commutator (ETC) is calculated in the usual formalism. (J. M. Cornwall and R. Jackiw, Phys. Rev. D4, 367 (1971).) Alternatively we can calculate $\{\psi^+(\vec{x},t),\ \psi(0)\}$ in terms of the ETC by Taylor expanding $\psi^+(\vec{x},t)$ in terms of $\frac{\partial^n \psi^+(\vec{x},0)}{\partial t^n}$, constructed using the Hamiltonian; the fact that the leading singularity does not depend on masses or on (pseudo) scalar interactions follows immediately from dimensional considerations. (D. J. Gross and S. Treiman, Phys. Rev. D4, 1059 (1971), C. H. Llewellyn Smith, Phys. Rev. D4, 2392 (1971).)

5. For a review and reference see S. L. Adler in "Lectures in Particles and Quantum Field Theory", Vol. 1 (MIT Press, Cambridge, Mass., and London, England (1970)).

6. The relations which obtain in the quark model have been reviewed by C. H. Llewellyn Smith SLAC-PUB-958 (1971) (to be published in Physics Reports).

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8. This is discussed in some detail by R. L. Jaffe, Phys. Letters 37B, 517 (1971) and SLAC-PUB-999 (to be published in Phys. Rev.) and J. C. Polkinghorne,Nuovo Cimento 8A, 592 (1972).

9. J. D. Bjorken and E. A. Paschos, Phys. Rev. D1, 1450 (1970) and references therein.

S. D. Drell and T. M. Yan, Annals of Physics 66, 578 (1971) and references therein.

P. V. Landshoff and J. C. Polkinghorne, paper contributed to this conference and references therein.

S. Brodsky, F. E. Close and J. Gunion, SLAC-PUB-1012 (to be published in Phys. Rev.).

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11. R. L. Jaffe (Ref. 8).

12. R. Brandt and G. Preparata, Brookhaven preprint BNL 16183 (1971).

13. R. L. Jaffe, private communication.

14. S. M. Berman, private communication.

15. If the proton's momentum is symmetrically distributed over the partons: $\langle x \rangle = \langle \frac{1}{N} \rangle$, where N is the number of partons. Measurements as $x \to 1$ therefore probe small N configurations.

16. C. G. Callan and D. J. Gross, Phys. Rev. Letters 25, 156 (1969).

17. For a review of the data see the accompanying talk presented by W. Toner at this conference.

18. In lines 5 and 6 we could obtain equalities for ϵ if we had sufficient neutrino data. The numbers in the table are based on the total neutrino cross section. (I. Budagov et al., Phys. Letters 30B, 364 (1969) and G. Myatt and D. H. Perkins, Phys. Letters 34B, 542 (1971). The values adopted for the integrals over the SLAC-MIT data are those used in Ref. 6; note that there is an unknown systematic error as the data were extrapolated to $\omega = \infty$ (new evaluations of these integrals will be available shortly after final data analysis--H. Kendall, private communication).

19. In calculating ϵ we have assumed that chiral symmetry obtains on the LC, i.e., that $F_2(VV) = F_2(AA)$. (Note that failures of chiral symmetry can be bounded since $\frac{|\sigma^{\nu} - \sigma^{\bar{\nu}}|}{\sigma^{\nu} + \sigma^{\bar{\nu}}} \leq \frac{2\gamma}{1+\gamma^2}$ for isocalar targets if $|V| = \gamma |A|$). Most of the results for ϵ can be obtained using results in a paper by Nachtmann (Phys. Rev. D5, 686, 1972). The essential ingredient in the derivations is that the distributions U^I_{Iz} for partons with isospin (I, Iz) are not independent; averaging over neutron and proton targets — $U^I_{Iz}(p) + U^I_{Iz}(n) = U^I_{Iz}(p) + U^I_{-Iz}(p)$ is isotropic i.e., independent of I_z. (One can think of the u_i's as probabilities for nucleon $\rightarrow$ i + x and note that the states x can be summed incoherently; introducing (positive semi-definite) reduced U's for each I value of x this leads immediately to Nachtmann's inequalities $2U_p(p) \geq U_n(p)$, $2U_{\bar{n}}(p) \geq U_{\bar{n}}(p)$ in the quark model.)

20. The references for the different bounds are: Quark model; C. H. Llewellyn Smith, Nucl. Phys. B17, 277 (1970); Chiral Symmetry; J. D. Bjorken and E. A. Paschos, Phys. Rev. D1, 3151 (1970).

Sakata/Fermi Yang models; C. H. Llewellyn Smith, SLAC-PUB-958 (to be published in Physics Reports); Q integral, $I \lesssim 1/2$; O. Nachtmann, Phys. Rev. D5, 686 (1972).

21. D. P. Majundar, Phys. Rev. D3, 2869 (1971), O. Nachtmann, Orsay Preprint LPTHE 71/12 and Nuclear Physics B38, 397 (1972).

22. H. J. Lipkin, Phys. Rev. Letters 28, 63 (1972).

23. R. Budny, T. H. Chang and D. K. Choudhury, Oxford preprint 5/72.

24. The inequality

$$\frac{\sigma^{\nu p+\nu n}}{2} + \frac{\sigma^{\bar{\nu} p+\bar{\nu} n}}{2} = \frac{G^2 MER}{\pi} \int F_2^{ep+en} dx \lesssim (0.67 \pm 0.09) \frac{G^2 ME}{\pi}$$

might prove fatal. Note that R appears to be near its upper limit which would imply that the isoscalar contribution to F_2^e is 10% of the isovector; this is not very surprising since this is approximately the ratio observed in photoproduction.

25. Consider models in which the nucleon consists of three "valence quarks" (with a common momentum distribution V(x)) and an SU(3) symmetric Q-$\bar{Q}$ sea (with a distribution S(x)). Such models have been considered by Bjorken and Paschos, Landshoff and Polkinghorne, Kuti and Weisskopf and many others. These authors made models for S(x) and V(x). Here we express S and V in terms of F_2^{ep} and F_2^{en} and then use them to calculate σ^ν; in so far as other models fit the data their results must be the same as ours. (Note that these models are unlikely to be strictly correct since they imply that $\int (F_2^{ep} - F_2^{en}) \frac{dx}{x} = \frac{1}{3}$ and that $\frac{F_2^{en}}{F_2^{ep}} \gtrsim \frac{2}{3}$. In the region where this ratio is less than $\frac{2}{3}$ our fitted S(x)

is negative! However this simple model might provide a first approximation.) The model gives:

$\underline{\Delta S = 0}$

$$\frac{\sigma^{\nu p+\nu n}}{2} = (0.44\pm0.16)\frac{G^2ME}{\pi}$$

$$\frac{\sigma^{\bar{\nu}p+\bar{\nu}n}}{2} = (0.20\pm0.08)\frac{G^2ME}{\pi}$$

$$\frac{\sigma^{\nu n}}{\sigma^{\nu p}} = 1.7\pm0.6, \qquad \frac{\sigma^{\bar{\nu}n}}{\sigma^{\bar{\nu}p}} = 1.5\pm\begin{matrix}3.5\\0.5\end{matrix}$$

$\underline{\Delta S = 1}$

$$\sigma^{\bar{\nu}\,\substack{n\\p}} = \tan^2\theta_c\, \sigma^{\bar{\nu}\,\substack{n\\p}}_{\Delta S=0}$$

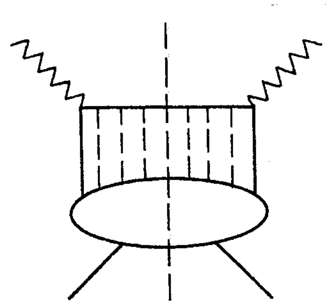
$$\sigma^{\nu p} = \sigma^{\nu n} = \tan^2\theta_c \frac{G^2ME}{\pi}\int(6F_2^{en} - 4F_2^{ep})dx$$

$$= \tan^2\theta_c(0.06\pm0.20)\frac{G^2ME}{\pi}\,.$$

The value for $\sigma^{\nu p+\nu n}$ is close to the CERN result: $\frac{\sigma^{\nu p+\nu n}}{2} = (0.52\pm0.13)\frac{G^2ME}{\pi}$.

26. S. D. Drell and T. M. Yan, Ref. 9.

27. If there is a vector interaction it is no longer true that there is no final state interaction; the quark propagates in the vector field (----) with the eikonal phase $\exp\left(ig\int_0^y \phi_\mu(z)dz^\mu\right)$ thus



However, the bilocal can still be factorized.

28. Ken Johnson (SLAC-PUB-1034) has shown that models in which quarks "stay in" have a remarkable self consistency.

29. D. G. Sutherland, Nucl. Phys. B2, 433 (1967).

M. Veltman, Proc. Roy. Soc. A301, 107 (1967).

30. J. S. Bell and R. Jackiw, Nuovo Cimento 60A, 47 (1969).

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35. R. J. Crewther, Cornell preprint CLNS-178 (1972); this paper contains a discussion of the analogue of the Bell-Jackiw-Adler anomaly for the divergence of the scale current, which has also been investigated by M. Chanowitz and J. Ellis SLAC-PUB-1028.

36. There are other results which obtain in parton models in which quarks are produced and might be abstracted by the courageous, e.g., for $e^+e^- \to \gamma \to$ hadron $+ \cdots\cdot$, quark parton models give $d\sigma(\pi^o) = d\sigma(\pi^\pm)$, $4 \gtrsim \frac{d\sigma(n)}{d\sigma(p)} \gtrsim \frac{1}{4}$, $4 \gtrsim \frac{d\sigma(K^o)}{d\sigma(K^+)} \gtrsim \frac{1}{4}$ etc. These results also follow from an operator expansion treatment of this process (J. Ellis, Phys. Letters 35B, 537 (1971)) (which need not imply the production of quarks) if the operators with the lowest dimensions in the expansion of $J_\mu(y)J_\nu(0)$ have $I \lesssim 1$, as is the case in the quark model (these considerations emerged from discussions with John Ellis).

37. Or Han-Nambu quarks. With paraquarks the usual quark model predictions for eN and νN scattering are unchanged but the prediction for pp $\rightarrow \mu\bar{\mu}$ + $\cdot\cdot$ decreases by a factor of 3 (R. L. Kingsley and C. Nash, Cambridge preprint 72/17).

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46. J. D. Bjorken in Proc. 1971 Symposium on Electron and Photon Interactions at High Energies (Lab. of Nuclear Studies, Cornell, 1971).

47. See B. L. Ioffe, Ref. 1.

48. J. Kogut, "Electroproduction and the "Size" of a Virtual Photon", Inst. of Advanced Study Preprint, 1971.

49. H. T. Nieh, Phys. Letters 38B, 100 (1972).

50. At nonasymptotic energies the slope may change because the relative importance of P and f exchanges may change with Q^2. Therefore it is necessary to know the behavior as a function of S and Q^2 in order to know whether R_γ has changed (Y. Avni, private communication).

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52. Models can easily be constructed in which the dip moves much less than in Harari's "minimalistic model" (Y. Avni, private communication).

53. For a review of the status of these theories in Jan. 1972 see: B. W. Lee, NAL-THY-34. Theories of this type were first considered by Salam and Ward (Ref. 57) and (independently) by Weinberg (Ref. 58); both Weinberg (Ref. 58) and Salam (Ref. 59) suggested that these theories might be renormalizable. The relevance of gauge theories has been repeatedly emphasized by Veltman in recent years.

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64. The Goldstone theorem requires manifest covariance. Like quantum electrodynamics, Yang-Mills theories can be written in forms which are not manifestly covariant. In manifestly covariant formulations there are Goldstone bosons but they decouple (e.g., in the R gauge the numerators of the propagator is $g_{\mu\nu} - \frac{k_\mu k_\nu}{k^2}$, but the "ghost" pole at $k^2 = 0$ is totally decoupled because the currents are conserved).
The Goldstone theorem and the Higgs phenomenon have been reviewed by G. S. Guralnick, C. R. Hagen and T. W. B. Kibble in Advances in Particle Physics Vol. 2 (Interscience, 1968).

65. See Refs. 62, 63 and 64. See G. Rajesakaran ("Yang-Mills fields and Theory of Weak Interactions" Tata Institute, Bombay, preprint, 1971) for a lucid introductory review.

66. D. J. Gross and R. Jackiw, MIT preprint 262 (1972).

67. C. Bouchiat, J. Iliopoulos and Ph. Meyer, Phys. Letters 38B, 519 (1972).

68. W. A. Bardeen has recently shown this using a modification of the techniques developed by 't Hooft and Veltman (Ref. 54) (private communication; CERN preprint in preparation).

69. S. L. Glashow has made a systematic investigation of the circumstances in which this occurs (private communication; Harvard preprint in preparation).

70. B. W. Lee, Ref. 53.

71. The way this cancellation occurs in Weinberg's model is explicitly exhibited in Ref. 73. See T. Appelquist and H. Quinn (Harvard preprint: "Divergence Cancellations in a Simplified Weak Interaction Model"; to be published) for a simplified model; this model suggests that the demand that WW scattering be well behaved at high energies would require the introduction of scalar ("Higgs") mesons (just as $e\bar{e}(\nu\bar{\nu}) \rightarrow W^+W^-$ motivates the introduction of Z's or heavy leptons).

72. This argument (which seems to be part of the folklore) was impressed on me by J. D. Bjorken.

73. S. Weinberg, Phys. Rev. Letters 27, 1688 (1971).

74. S. Weinberg, MIT preprint 246 (1971) (to be published).

75. H. H. Chen and B. W. Lee, "Experimental tests of Weinberg's theory of Leptons", U.C. Irvine preprint UCI-10P19-63 (1971) (to be published in Phys. Rev.).

76. The result depends on the mixing angle defined by $\tan\theta = g'/g$. This is fixed in another version of Weinberg's model (S. Weinberg MIT preprint 253 (1971); to be published).

77. S. L. Glashow, J. Iliopoulos and L. Maiani, Phys. Rev. D2, 1285 (1970).

78. G. Myatt (private communication).

79. R. Budny (private communication); this result is for the "anomaly free" version of the model (Ref. 67). The parton model is that used by Budny in his Oxford preprint 6/72.

80. H. Georgi and S. L. Glashow, "Unified Weak and Electromagnetic Interactions without Neutral Currents", Harvard preprint (1972).

81. J. D. Bjorken, unpublished.

82. T. Appelquist and H. Quinn (private communication; Harvard preprint in preparation).

83. For tests of the symmetry properties of currents see A. Pais, Phys. Rev. D5, 1170 (1972) and the references cited there.