NEUTRINO REACTIONS

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Abstract

Neutrino reactions are reviewed from a theoretical point of view.

Contents

Foreword

First results from high statistics neutrino experiments underway at the CERN SPS appeared in preprint form just after the 1977 CERN school. Not only do the data have much smaller errors than data from previous experiments but they differ qualitatively in some cases. The experimental situation is evolving extremely rapidly. A considerable amount of additional SPS data has already been shown in seminars and summer schools and much more will be available in print by the time the proceedings of the CERN school are published. It is therefore not sensible for me to review the data now available in detail or to attempt to squeeze all possible information from them. Instead, I have concentrated on the ideas which underlie the theoretical interpretation of neutrino experiments. I have not shown any SPS data, although some of the results are quoted in the text (and I have, of course, taken cognisance of all the data I have seen in deciding what to write). The reader is referred to the proceedings of the forthcoming Hamburg Lepton-Photon Symposium for authoritative reviews of the experimental situation at the end of the summer of 1977.

Refs.1-11 are to some of the many previous reviews of various aspects of neutrino physics which the reader may find helpful.

1. Introduction

1.1 Why Do Neutrino Experiments?

Most of our knowledge of Weak Interactions derives from:

a) Study of the decays of the eight known stable non-charmed particles (μ, π, K, Λ, Σ, Ξ, Ω) and the charmed D meson (not much is known yet about charmed baryons).

b) Study of mu-capture $\mu^-(A, Z) \to \nu(A, Z-1)$.

c) Study of parity violation in nuclei and atoms.

d) Neutrino reactions.

In a-c the subject matter is limited and the momentum transfers involved are small. Neutrino experiments allow us to study additional processes in a new domain. By doing neutrino experiments we might hope

a) to learn the domain of validity of the "conventional theory" of weak interactions (which is necessarily wrong, as discussed below);

b) to find new phenomena which will help us to establish a new theory (presumably a gauge theory);

c) given a theory, to probe the structure of hadrons.

1.2 The Conventional Theory of Weak Interactions

An essential prerequisite for studying neutrino reactions is a knowledge of the elements of conventional weak interaction theory and its limitations. The brief remarks which follow are intended to establish notation and give the reader an idea of some things he/she ought to know and where they can be learned[12]. In the conventional theory, weak interactions were described by an "effective Lagrangian":

$$\mathcal{L}_{eff.} = \frac{G}{\sqrt{2}} J^{+}_{\lambda}(x) J^{\lambda}(x)$$

where $G = \frac{10^{-5}}{M_p^2}$ is Fermi's constant (in the "natural" units $\hbar=c=1$ used here). The current J_λ is given by

$$\begin{aligned} J_\lambda(x) = {} & \bar{\nu}_\mu(x)\gamma_\lambda(1-\gamma_5)\mu(x) \\ & + \bar{\nu}_e(x)\gamma_\lambda(1-\gamma_5)e(x) \\ & + \bar{u}(x)\gamma_\lambda(1-\gamma_5)[d(x)\cos\theta_c + s(x)\sin\theta_c] \\ & + \cdots\cdots \end{aligned}$$

where $\nu_\mu(x)$ is the field operator of the mu-neutrino, u(x) the field of the up quark etc., θ_c is the Cabibbo angle and the γ matrices are those used in Bjorken and Drell. We now know that there is an extra charm changing piece of the current, which couples the charmed field c(x) predominantly to s(x), all the data being consistent with the Glashow, Iliopoulos and Maiani (GIM) form:

$$\bar{c}(x)\gamma_\lambda(1-\gamma_5)[s(x)\cos\theta_c - d(x)\sin\theta_c].$$

We shall see that neutrino experiments are a powerful way to search for additional weak currents. $\mathcal{L}_{eff}$ above only describes charged current interactions. We shall discuss the effective Lagrangian which describes neutral current interactions (discovered in 1973) in Section 9.

The effective Lagrangian works very well for leptonic and semi-leptonic charged current processes in Born approximation, i.e. its matrix elements represent the measured matrix elements accurately in terms of a few parameters. For example, all features of semi-leptonic baryon decays are well described in terms of G, θ_c and the F/D ratio for the axial current matrix elements[13]. However:

a) $\mathcal{L}_{eff}$ embodies many untested features. A few examples:

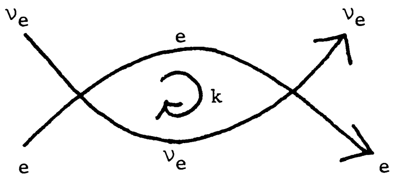
- it predicts the existence of the process $\nu_e e \to \nu_e e$.
- it forbids (e.g.) $\Delta S=4$ interactions.

- it prohibits the existence of "second class currents". At the moment the limits are very weak and in fact there is some evidence that second class currents exist[14]. Further experiments on this question are in progress and their outcome is eagerly awaited since second class currents are <u>anathema</u> for gauge theories[15].

b) $\mathcal{L}_{eff}$ has some trouble in accounting for the large enhancement of $\Delta I=1/2$ amplitudes and the depression of $\Delta I=3/2$ amplitudes, relative to simple minded estimates, which are observed in all $\Delta S=\pm 1$, $\Delta C=0$ non-leptonic decays[16]. It is conceivable that dynamical effects in the matrix elements may account for this, but it may also require a modification of the Lagrangian.

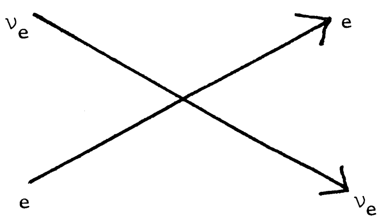
c) The theory must be wrong!

If treated as a real Lagrangian, $\mathcal{L}_{eff}$ leads to non-sense in higher orders, e.g. it gives a diagram



with an amplitude which diverges violently ($\sim \int \frac{d^4k}{k^2}$) when the integral over the unobserved momentum flowing around the loop is carried out. In contrast to the situation in Quantum Electrodynamics, this divergence cannot be absorbed by redefining the parameters in the original Lagrangian. It can only be cured at the price of introducing new (arbitrary) parameters. More and more are needed in each order rendering the theory devoid of all predictive power.

We cannot simply forget higher orders, however, since the lowest order taken on its own produces non-sense at high energy, e.g. corresponding to the diagram

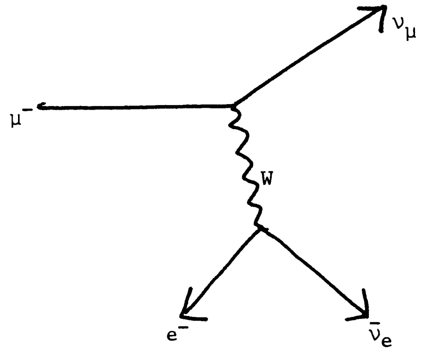


there is a cross-section $\sigma \sim G^2 s$ (a result which is obvious on dimensional grounds since we can safely set $m_e=0$ at high energy), but the amplitude is purely s wave so eventually this conflicts with the unitarity bound $\sigma <$ const./s.

The source of these troubles is the fact that G has dimensions M^{-2} so it would seem that the introduction of massive intermediate vector bosons $W^\pm$ might effect a cure. We write

$$\mathcal{L} = g\left(W_\lambda J^\lambda + h.c.\right)$$

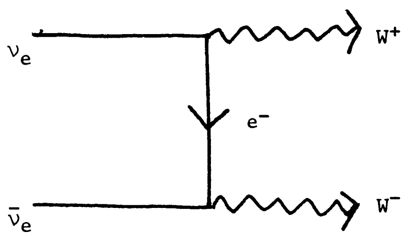
so that (e.g.) mu-decay is now described by the diagram



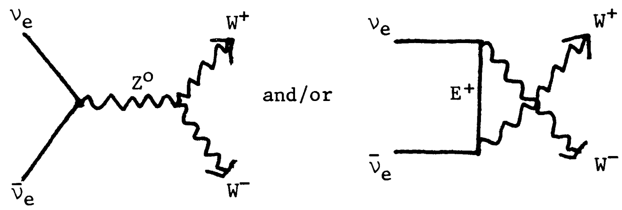
Since it is known (see Section 10) that $M_W>8.4$ GeV, we can safely neglect Q^2 in the propagator in this process and identify

$$\frac{g^2}{M_W^2} = \frac{G}{\sqrt{2}} .$$

Although g is dimensionless (like e in QED), this theory still leads to violent non-renormalizable divergences and gross violations of unitarity in processes involving longitudinal W's, whose polarization vectors ε_λ behave like k_λ/M_W (where k_λ is their four momentum) at high energies. For example, corresponding to the diagram



there is a cross-section $\sigma \sim g^4 s/M_W^4$. In a theory which can be treated perturbatively, this can only be cured by cancellation with other contributions of the same order[17], the only possibilities being neutral current or heavy lepton exchange:



In fact it is possible to construct unitary renormalizable gauge theories of weak and electromagnetic interactions with just these features.

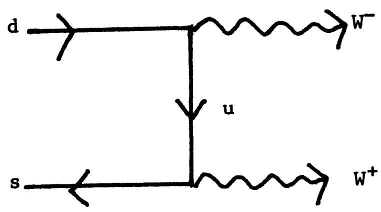
These theoretical arguments lead us

a) to expect the failure of $\mathcal{L}_{eff}$ at very high energy.

b) to think that the cure will involve W's (with $g\sim e$, we find $M_W\sim 100$ GeV).

c) to have anticipated the existence of neutral currents and heavy leptons.

d) to have anticipated the existence of charm (the contribution of u quark exchange to $d\bar{s} \to WW$



is sick, like the electron exchange contribution to $\nu_e\bar{\nu}_e \to W^+W^-$; however, the worst behaved term, which is independent of the quark mass, is exactly cancelled by the c quark contribution).

All these ideas can be tested in neutrino experiments. Actually, there are other conceivable cures for the diseases of $\mathcal{L}_{eff}$. In all cases charm and the GIM mechanism are probably needed to render $K^0_L \to \mu^+\mu^-$ and $M_{K^0_L} - M_{K^0_S}$ small[6]. In most cases substantial neutral current effects would probably be induced by higher order effects (but a pure V-A structure would then presumably be expected). Nevertheless, gauge theories seem to me the only serious approach to the construction of weak interaction models at present because of their elegance (and also because of the quantitative success of the SU(2)xU(1) model in describing neutral current data which will be discussed in Section 9). Much of the theoretical discussion in these lectures will be in the framework of gauge theories, which is now the standard conventional framework. In fact the remarks here are intended to direct any reader who is not reasonably versed in gauge models to the abundant pedagogical material now available[12].

1.3 Orders of Magnitude

The cross-section for the process $\nu n \to \mu^- p$ involves a form factor which limits the momentum transfer to $Q^2 < M_p^2$ so that

$$\sigma_{\nu n \to \mu^- p} \sim G^2 M_p^2 \sim 10^{-38}\ \mathrm{cm}^2$$

With this cross-section, the mean free path of a neutrino in lead would be one astronomical unit (the diameter of the earth's orbit around the sun). To first approximation, inelastic neutrino reactions are "scale invariant" (see Section 3) so that $\sigma \sim G^2 s$. In fact

$$\sigma^{\nu} \sim 0.8\, E_\nu \times 10^{-38}\ \mathrm{cm}^2\ \mathrm{GeV}^{-1}$$

$$\sigma^{\bar{\nu}} \sim 0.3\, E_{\bar{\nu}} \times 10^{-38}\ \mathrm{cm}^2\ \mathrm{GeV}^{-1}$$

We should now consider how experiments can be done with these very small cross-sections.

1.4 Experiments

These lectures are about the theoretical interpretation of neutrino data. This requires some idea of how the experiments are done. A very clear discussion of neutrino beams and of bubble chamber and counter experiments can be found in Steinberger's lectures at the 1976 CERN school[9], which the uninitiated are advised to read. I will not repeat this material in these written notes. However, it is worth mentioning neutrino emulsion experiments which make it possible to look for very short lived particles leaving tracks or gaps of a few microns close to the primary vertex. Large amounts of emulsion (e.g. 20 litres or more) are needed which makes it essential to employ ancilliary devices which can locate the position of interesting vertices before scanning. For example, in the successful experiment already carried out at FNAL[18] wire chambers etc. were located downstream. At CERN an emulsion stack will be placed in front of the bubble chamber BEBC so that the photographs and information from the EMI (external muon identifier) can be used to select interesting events (e.g. multimuon events, $^-e^+K^0_S$ events or very high energy events) and locate their vertices. In the lectures I also mentioned the possibility of undersea detection of the interaction of cosmic ray neutrinos using a large ($\sim$ 1 Km^3!) array of photomultipliers to detect Cerenkov radiation; the interested reader is referred to the literature on this project[19].

1.5 Survey of Neutrino Reactions

The following brief notes provide an overview of the subject and some references to sources of further information on topics which will not be discussed here.

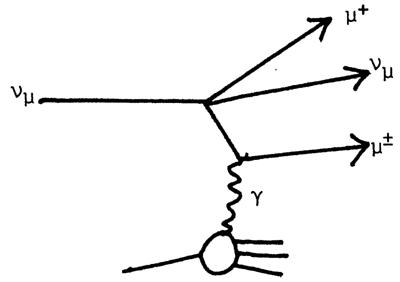
a) Leptonic processes:

Could study: $\nu_\mu e \to \mu\nu e$ (inverse μ decay[1])

$\nu_\mu e \to \nu_\mu e$ (neutral current: see Section 9).

$\sigma \sim G^2 s = 2m_e G^2 E_\nu \sim 10^{-41} E_\nu \mathrm{cm}^2 \mathrm{GeV}^{-1}$ ➔ experiments hard.

Was hoped[1] to study



but the "background" from $\nu_\mu N \to \mu^- \mu^+$ events (due to charm - see Section 5) makes this impossible.

b) Semileptonic processes

Charged currents

Exclusive:

$\nu n \to \mu^- p$ $\nu N \to \mu\Delta$	Form factors[1] Selection rules[1] e.g. $\Delta I=1 \quad \to \quad \dfrac{\nu p \to \mu\Delta^{++}}{\nu n \to \mu\Delta^{+}} = 3$
$\bar{\nu}p \to \mu^+\Lambda$ $\nu p \to \mu^- C$	Test Cabibbo at large Q^2 [1] Investigate charmed particles[20]

Inclusive:

$\nu N \to \mu x$	Test ideas of nucleon structure (partons, asymptotic freedom, etc.). See "new" pieces of current - changes in σ, $\frac{d\sigma}{dy}$ etc. (See Section 8).

Neutral Currents

$\nu p \to \nu p$ $\nu p \to \nu\Delta$ $\nu N \to \nu x$, etc.	Find space-time, I-spin, etc. properties of N.C. (see Section 9).

c) New Phenomena

Found: $\nu N \to \mu^- \mu^+ X$ - due to charm production? See Section 5.

Might be found: $\nu N \to M^+ X$ - See Section 8.
$\quad \hookrightarrow \mu^+ \ldots$

Expected at very high energy: $\nu_\mu N \to \mu^- W^+ X$ - See Section 10.

d) Solar Neutrinos - the reasons why fewer neutrinos seem to arrive at the Earth than expected, according to the standard solar model, is outside the scope of these lectures[21].

e) Neutrino Oscillations[22]

Suppose lepton number is not conserved and $m_\nu \neq 0$ (as occurs in several fashionable models). Then in general the states ν_μ and ν_e which are coupled to μ and e by the weak current will not coincide with the eigenstates ν_1 and ν_2 of the mass matrix

$$\begin{pmatrix} m_\nu^\mu & \tilde{m} \\ \tilde{m} & m_\nu^e \end{pmatrix}$$

in analogy to the situation with K^0, $\bar{K}^0$ and K_1, K_2 (if there are more than two neutrinos they will also get mixed). Then a ν_μ produced, e.g. in $\pi \to \mu\nu_\mu$ decay will evolve in time as

$$\nu_1 \cos\xi \, e^{-iE_1 t} + \nu_2 \sin\xi \, e^{-iE_2 t}$$
$$= a(t)\,\nu_\mu + b(t)\,\nu_e$$

with

$$|b|^2 = 1 - |a|^2 = 4\cos^2\xi \sin^2\xi \sin^2\left(\frac{\Delta E t}{2}\right).$$

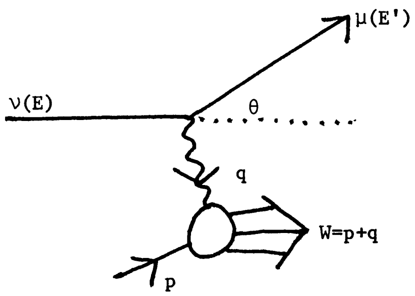
The wavelength of the oscillations is

$$\lambda = \frac{5p}{\Delta} \quad \text{metres}$$

where p is the momentum in MeV (assumed large compared to $m_{\nu 1}$ and $m_{\nu 2}$) and $\Delta = |m_{\nu 1}^2 - m_{\nu 2}^2|$ in eV^2. The oscillations allow supposedly pure ν_μ beams to produce e^- and induce a sterile $\bar{\nu}_\mu$ component in $\bar{\nu}_e$ beams from reactors with energy $< m_\mu$ (with enough neutrinos, the solar neutrino experiment could be explained in this way). Reactors are best suited to look for oscillations if $\Delta \ll 1$ and accelerators if $\Delta \gg 1$. Present limits are displayed in Fig.1.

2. Kinematics

A detailed formal discussion of kinematics can be found (e.g.) in reference 1. We sketch the main results here. Energies and momenta are defined by



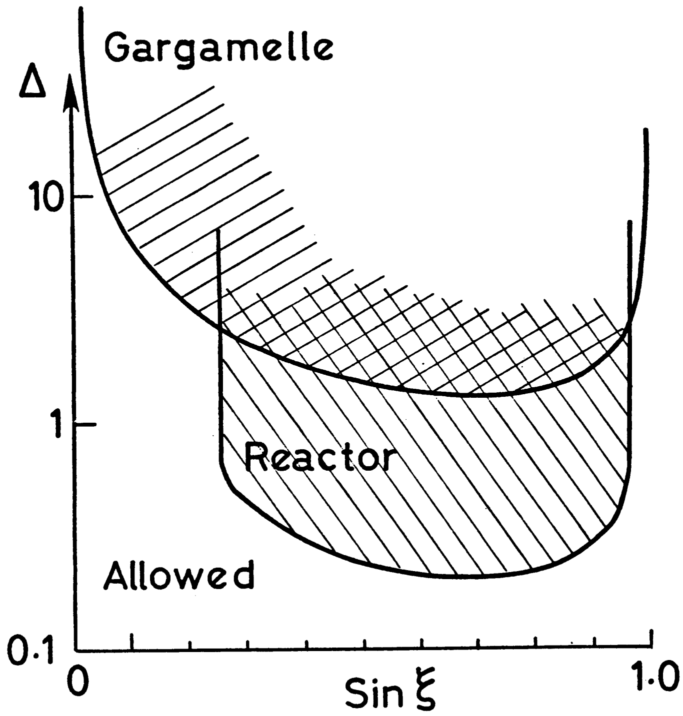


Fig.1. 90% confidence level constraints on Δ (in eV^2) and $\sin\xi$. The Gargamelle data[23] give $|b|^2 < 0.003$ with $\langle p_\nu \rangle = 1.5$ GeV and I have assumed an average distance from the neutrino source to the detector of about 50m. The reactor experiment[24] gives $|b|^2 < 0.25$ for $\langle p_\nu \rangle = 3$ MeV at a distance of 6m.

and we shall use the variables:

$$Q^2 = -q^2 = \vec{q}^{\,2} - q_0^2 \geq 0$$
$$\nu = q \cdot p = M(E - E')$$
$$x = Q^2/2\nu$$
$$y = \frac{q \cdot P}{ME} \simeq \frac{2\nu}{s}$$
$$W^2 = (p+q)^2$$
$$= M^2 + 2\nu - Q^2$$
$$s = (p + k_\nu)^2$$
$$= M^2 + 2p \cdot k_\nu \,.$$

The allowed kinematic regime is bounded by a triangle in the Q^2-ν plane, as shown in Fig.2.

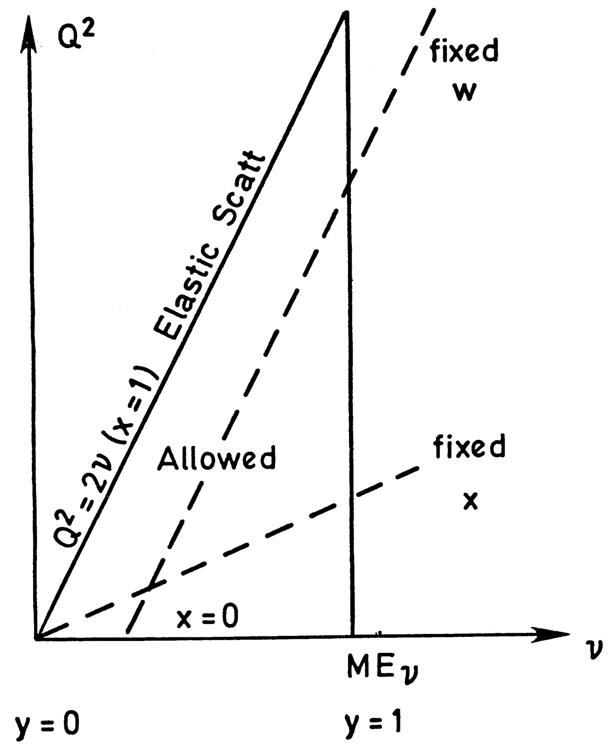


Fig.2. The allowed region in the Q^2-ν plane.

Those not familiar with this plot are urged to work out the boundaries, etc.

The amplitude corresponding to the Feynman diagram above is

$$\frac{g^2 j^\mu(0)\left(-g_{\mu\nu}+\frac{q_\mu q_\nu}{M_W^2}\right)J^\nu(0)}{q^2-M_W^2}$$

where the leptonic current is known

$$j_\mu(x) = \bar{\mu}(x)\gamma_\mu(1-\gamma_5)\nu_\mu(x)$$

and

$$J_\nu(x) = \langle F|\,J_\nu^{Weak}(x)\,|p\rangle.$$

We set $m_\mu=0$ in which case

1) $\frac{\partial j_\mu(x)}{\partial x_\mu}=0$ so that j_μ has only three independent components and the current/virtual W has J=1.

2) The $\mu^-(\mu^+)$ is purely left (right) handed.

There is then no lepton spin sum involved and the current/virtual W is in a pure state. Explicit calculation gives

$$j_\lambda \sim \sqrt{\frac{E}{2E'}}\,\epsilon_\lambda^L+\sqrt{\frac{E'}{2E}}\,\epsilon_\lambda^R+\epsilon_\lambda^0$$

where we have assumed $\nu^2 \gg M^2Q^2$, and

$$q = (q_0,0,0,|\vec{q}|)$$
$$\epsilon_\lambda^{L,R} = \frac{1}{\sqrt{2}}(0,1,\pm i,0)$$
$$\epsilon_\lambda^0 = \frac{1}{\sqrt{Q^2}}(|\vec{q}|,0,0,q_0).$$

The cross-section has the form

$$d\sigma \sim |j_\lambda J^\lambda|^2$$

for some fixed final state. If the hadrons are rotated through an azimuthal angle ϕ about $\vec{q}$ (with the leptons fixed), the components of J_λ change thus: $J_L\to e^{-i\phi}J_L$, $J_R\to e^{+i\phi}J_R$, $J_0\to J_0$ in the helicity basis above. An integration over ϕ removes the interference terms so that the neutrino cross-section can be expressed in terms of the cross-sections for helicity +1, -1 and 0 (R, L, 0) currents/virtual W's:

$$\frac{d\sigma}{dQ^2 d\nu d\Gamma} = \frac{(\nu-Q^2/2)}{\pi}\,\frac{G^2Q^2E'}{2\pi\nu^2E} \times\left[\frac{d\sigma^0}{d\Gamma}+\frac{E'}{2E}\frac{d\sigma^R}{d\Gamma}+\frac{E}{2E'}\frac{d\sigma^L}{d\Gamma}\right]$$

where Γ represents all the variables which characterise the final hadronic state (apart from ϕ) and we have assumed $Q^2 \ll M_W^2$. For $\bar{\nu}$ the expression has the same form with

$$d\sigma^{\nu}_{0,R,L} \to d\sigma^{\bar{\nu}}_{0,L,R}$$

($d\sigma_i^\nu \neq d\sigma_i^{\bar{\nu}}$ in general since one is a W^+ and the other a W^- cross-section. Integrating over Γ, we obtain (still for $\nu^2 \gg M^2Q^2$ - see ref. 1 for the exact formula)

$$\frac{d\sigma}{dx\,dy} = \frac{(\nu-Q^2/2)}{\pi}\,\frac{G^2xME}{\pi} \times\left[\sigma_L+(1-y)^2\sigma_R+(1-y)\sigma_0\right]$$
$$= \frac{G^2ME}{\pi}\left[(1-y)F_2+y^2xF_1-y(1-y/2)F_3\right]$$

where $F_{1,2,3}(\nu, Q^2)$ are linear combinations of $\sigma_{L,R,0}(\nu, Q^2)$ which are frequently employed (see e.g. ref. 1). Remarks:

1) Fixed values of ν and Q^2 correspond to different y for different E_ν, so the σ_i's/F_i's can be separated experimentally.

2) Assuming charge symmetry for the $\Delta S=\Delta C=0$ current and neglecting $\Delta S\neq 0$, $\Delta C\neq 0$ pieces, the W^+ and W^- cross-sections are related by

$$\sigma_i^{\nu n\to F} = \sigma_i^{\bar{\nu}p\to\bar{F}}$$
$$\sigma_i^{\nu p\to F} = \sigma_i^{\bar{\nu}n\to\bar{F}}$$

or equivalently

$$F_i^{\nu\binom{p}{n}\to F} = F_i^{\bar{\nu}\binom{n}{p}\to\bar{F}}$$

where $\bar{F}$ is obtained from F by a reflection in isospin space. For an I=0 target summing over F we therefore have

$$\sigma_i^\nu = \sigma_i^{\bar{\nu}}$$
$$F_i^\nu = F_i^{\bar{\nu}}$$

Note that this does not mean that $d\sigma^\nu = d\sigma^{\bar{\nu}}$.

3. Scale Invariance

Let us suppose that as ν, $Q^2 \to \infty$ all masses and parameters with dimensions can be neglected so that the σ_i depend only on the variables ν and Q^2. In this case the σ_i, which have dimensions $L^2 = M^{-2}$ in natural units with $\hbar = c = 1$, would have to behave as

$$\sigma_i(\nu, Q^2) \to \frac{1}{Q^2}\, g_i(x)$$

where $x \equiv \frac{Q^2}{2\nu}$ and likewise the dimensionless F_i would behave as

$$F_i(x, Q^2) \to F_i(x).$$

This postulate, known as Bjorken scaling, is far from obvious. In general it is certainly not true that masses etc. can be neglected at high energy and that cross-sections depend only on dynamical variables with dimensions. Such an idea would imply (e.g.) the wrong result that $\sigma_{pp}^{tot} \sim 1/s$; in the case of pp collisions the radius of the proton is another quantity with dimensions which is obviously relevant for the construction of σ^{tot}. However, the size of the virtual photon (current) in $ep \to e$ ... ($\nu p \to \mu$ ...) decreases like $1/\sqrt{Q^2}$ (technically, $\sigma_{\gamma^* p}$ is proportional to the forward $\gamma^* p \to \gamma^* p$ amplitude; if γ^* is absorbed at x_μ and emitted at y_μ, then deep inelastic scattering probes this amplitude in the limit $x_\mu - y_\mu \to 0$). Hence it might be thought that transverse dimensions play no role in $\sigma_{\gamma^* p}$ which might therefore satisfy Bjorken scaling. However, this idea is not exactly true either - at least in perturbation theory. As everyone who has calculated non-trivial Feynman diagrams knows, there are logarithmic divergences in the limit $m \to 0$. This gives rise to violations of Bjorken scaling (the g_i being functions of $\ln(Q^2/m^2)$, as well as x).

Experimentally, it was thought at one time that scaling might turn out to be exact. Some of the SLAC-MIT data[25)] which led to this conclusion are shown in Fig.3. Subsequently, improved data from the SLAC-MIT group[26)] and data at higher Q^2 from FNAL[27)] have shown that $F_2(x, Q^2)$ decreases with Q^2 at large x and increases with Q^2 at small x, as shown in Figs.4. Some - perhaps all - of the decrease at large x can be attributed to the wrong choice of scaling variable. For example, suppose

$$F_2(x, Q^2) = f\left(x' = \frac{x}{1 + \frac{M^2 x}{Q^2}}\right)$$

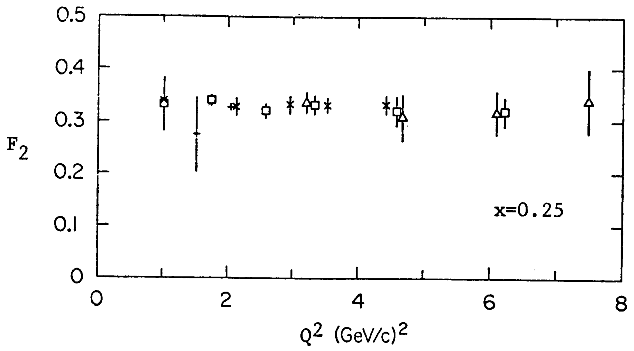


Fig.3. $F_2(x, Q^2)$ for ep scattering at x=0.25 as a function of Q^2 for W > 2 GeV[25)].

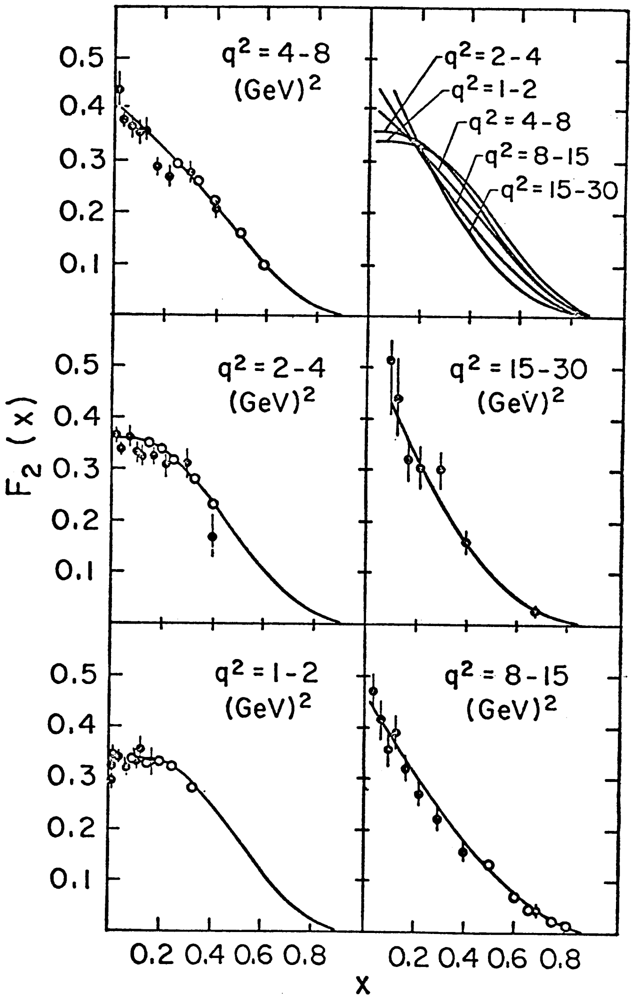


Fig.4. $F_2(x, Q^2)$ for ep scattering vs. x for various Q^2 ranges [27)]. ● FNAL data. ○ SLAC-MIT data. The lines, which are fits, are collected in the top right hand corner which shows the Q^2 dependence.

then F_2 would decrease with Q^2 at fixed x since f decreases with increasing x'. There are theoretical reasons[28)] to prefer some variable like x' and some of the decrease at large x is probably due to this effect. The rise at small x is also difficult to interpret as an unambiguous indication of scaling violation since some of it is probably due to the associated production of charmed particles (if scaling were exact, this contribution would also scale for $Q^2 \gg M_c^2$). However, if it were all due to charmed particle production this would surely have shown up in trimuon production by now. In fact I believe that the data show genuine scaling violation at both large and small x although data at higher Q^2 are needed to establish this completely.

In neutrino production, scale invariance would immediately predict

$$\sigma \propto G^2 s = 2MG^2 E_\nu^{lab}.$$

and

$$\langle Q^2 \rangle \propto s.$$

This is consistent with the Gargamelle data[29)] for E < 15 GeV (Fig.5). However, very recent SPS data

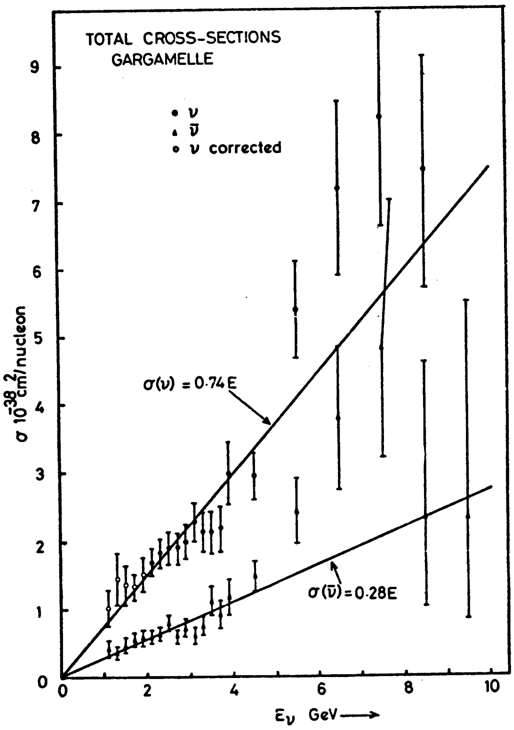


Fig.5. Total cross sections for neutrinos and anti-neutrinos[29].

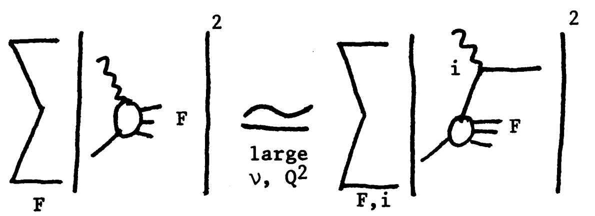
show that σ/s falls at higher energy[30] - despite the fact that there must be an increasing contribution from charm production. We discuss this clear indication of scaling violation further in Section 6.

Despite the experimental evidence that scaling may be violated, we shall now turn to the naive parton model which leads to exact scaling. As we shall see later, the model can easily be adapted to include scaling violations and, at the least, provides a very useful first approximation even in its naive form.

4. The Parton Model

4.1 Dogma

The basic idea of the parton model is that at large ν and Q^2 currents scatter incoherently off point-like constituents of the nucleon, called partons, which behave "as if" free during the interaction[31]. It is assumed that final state interactions can be neglected in calculating the total cross-section. Pictorially



The physics behind this dogma depends on the fact that the time τ which controls deep inelastic scattering is of order $1/\sqrt{Q^2}$. It might be argued that

a) As $Q^2\to\infty$, the strong interactions have no time to rearrange the constituents during the scattering processes and the current "sees" a frozen state of non-interacting partons. This proposition is false in renormalizable field theories, which leads to scaling violations, as will be discussed in Section 6.

b) The final state interactions, which must play an essential role in the formation of the final state if partons are confined quarks, act on a time scale $\tau' \gg \tau$. They may therefore be neglected in calculating the total cross-section. This proposition might be true, e.g. it is true in two-dimensional QCD, a (superrenormalizable) model with exact scaling and confinement in which the parton model works[32] and there is some truth in it in non-relativistic potential models[33]. A classical analogy may help: consider a particle attached to a point by an elastic band which is initially limp. When the particle is struck, it responds as a free particle. Later the band becomes taut and restrains it. However, this is irrelevant in calculating the total cross-section which is proportional to the probability that something happens. Once the particle moves off something has happened so the total cross-section is determined by the inital free particle response.

Undaunted by the fact that one part of the dogma fails in renormalizable field theories, let us move on to explore its consequences.

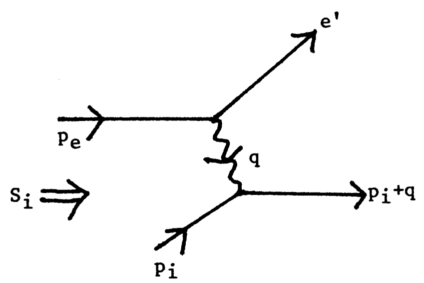
4.2 $ep\to eX$

Let the momentum p_i of the i^{th} parton be $\vec{p}_i = x\vec{p} + \vec{p}_T$, where $\vec{p}$ is the target momentum and $\vec{p}.\vec{p}_T = 0$. We will use the frame in which $q_0 = 0$, where $|\vec{p}| \simeq \nu/\sqrt{Q^2}\to\infty$ in the Bjorken limit ν, $Q^2\to\infty$ with $x = Q^2/2\nu$ fixed. According to the dogma, $|\vec{p}_T|$ will be fixed independent of ν and Q^2 in this limit. Thus $p^i_\mu \simeq x p_\mu$ in the Bjorken limit. The ep cross-section is then

$$\frac{d^2\sigma}{dQ^2 d\nu} = \sum_i \int dx\, f_i(x)\, \frac{d^2\sigma_i}{dQ^2 d\nu} \qquad (4.1)$$

where $f_i(x)$ is the probability of finding a parton of type i with fraction x of the momentum.

We now consider the cross-section $d\sigma_i$ for elastic e-parton scattering



which may be written

$$\frac{d\sigma_i}{dQ^2} = \frac{4\pi\alpha^2 Q_i^2}{Q^4}\, g_i(Q^2/s_i) \qquad (4.2)$$

where $Q_i|e|$ is the parton charge and we have assumed that Q^2 and $S_i = (p_e+p_i)^2 \simeq 2xp_e.p \simeq xs$ are large so that the parton mass can safely be neglected in Born approximation. Eqn.4.2 then follows from dimensional analysis. The function g_i, which depends on on the spin of the parton, is easy to calculate and is given by

$$\begin{array}{ll} \dfrac{1+(1-Q^2/S_i)^2}{2} & \text{for spin } \tfrac{1}{2} \\ (1-Q^2/S_i) & \text{for spin } 0. \end{array} \qquad (4.3)$$

The dependence on $Q^2/S_i = \frac{1-\cos\theta^{ei}_{c.m}}{2}$ will be explained later (note the normalization to the Rutherford cross-section as $Q^2 \to 0$ in both cases).

For elastic scattering $(q+p_i)^2 = p_i^2$ or $Q^2 = 2q.p_i \simeq 2x\nu$. Thus we may write

$$\frac{d\sigma_i}{dQ^2 d\nu} = \frac{d\sigma_i}{dQ^2}\,\delta\left(\nu - \frac{Q^2}{2x}\right) \qquad (4.4)$$

which may be substitued into eqn.4.1 to give

$$\frac{d^2\sigma}{dQ^2 d\nu} = \sum_i f_i(x)\,\frac{2x^2}{Q^2}\,\frac{d\sigma_i}{dQ^2} \qquad (4.5)$$

with $x = Q^2/2\nu$.

Comparison with the general form of the differential cross-section (which is obtained from $d\sigma^\nu$ above by putting $\sigma_R = \sigma_L = \sigma_T, F_3 = 0$ and making the substitution $G^2 \to \frac{8\pi^2\alpha^2}{Q^4}$) yields the parton model formulae for σ_0 and σ_T, or equivalently the dimensionless structure functions F_1 and F_2, in terms of Q_i^2 and g_i. Independent of the parton spin

$$F_2 = \sum_i Q_i^2\, x f_i(x) \qquad (4.6)$$

but F_1 is spin dependent.

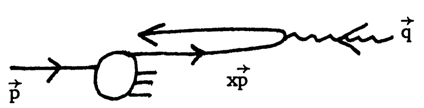
Remarks:

a) Since there are no parameters with dimensions in the parton model calculation, F_2 necessarily exhibits Bjorken scaling.

b) It follows from 4.3 that[34]

$$\frac{\sigma_0}{\sigma_T} = \infty \quad \text{for spin 0 partons}$$
$$= 0 \quad \text{for spin 1/2 partons}$$

This can be understood by helicity arguments since the parton and virtual photon meet head on in the frame we employ if we neglect p_T



Taking the z axis along $\vec{p}$, a spin zero parton has $J_z = 0$ initially and finally, and can therefore not absorb a virtual photon of helicity ± 1 so $\sigma_T = 0$, in this case. For a spin 1/2 parton, however, the bare electromagnetic vertex $\bar{u}\gamma_\mu u$ conserves helicity at high energy and $\Delta J_z^{\text{parton}} = 1$, which can only be supplied by a transverse photon, so that $\sigma_0 = 0$ in this case. The data, some of which are shown in Fig.6, show that $\frac{\sigma_0}{\sigma_T}$ is small and that this ratio could vanish as $Q^2 \to \infty$. We therefore assume that all electrically charged partons have spin 1/2 - presumably they are quarks.

c) The data (Fig.4) suggest that $F_2(x)$ is not zero in the limit $x \to 0$. It follows from eqn. 4.6 that $f_i(x) \sim 1/x$ at small x, i.e. partons have a Bremstrahlung spectrum as originally conjectured by Feynman[31] to explain σ_{pp}. Of course, the kinematic assumptions we have made are only justified

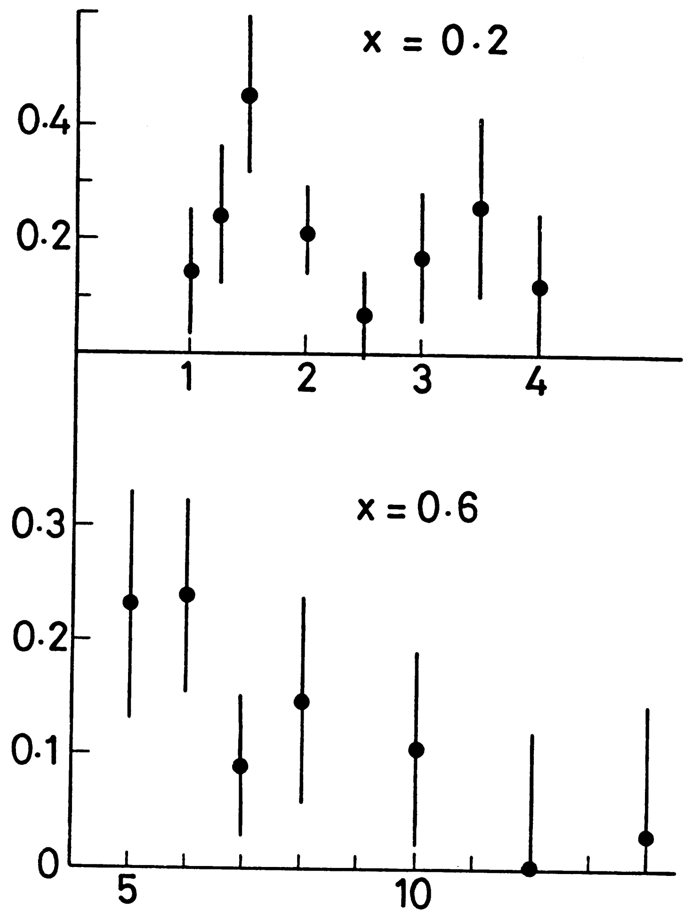


Fig.6. $\frac{\sigma_0}{\sigma_T}$ vs. Q^2 (in GeV^2) at two of the eleven values of $x = \frac{Q^2}{2\nu}$ for which data are given in ref. 35 and 26. The errors are statistical only (systematic errors decrease from about ± 0.13 at $Q^2 = 1$ to ± 0.09 at $Q^2 = 4$ for $x = 0.2$ and from ± 0.06 at $Q^2 = 5$ to ± 0.03 at $Q^2 = 14$ for $x = 0.6$).

for $xP \gg M_{\text{parton}}$ or $x \gg \frac{M^2_{\text{parton}}}{S_{\gamma^* p}}$. Integrating from this point we see that the parton multiplicity

$$\langle n_i \rangle = \int f_i(x)\,dx \sim \int \frac{dx}{x} \sim \ln S_{\gamma^* p}$$

grows with increasing energy.

d) Consider a configuration of N partons. If the protons' momentum were shared equally among them we would have

$$\langle x \rangle_N = \frac{1}{N}$$

In any case, we would expect the contribution to f(x) to be confined to small x for very large N and only to extend to large x for small N, i.e. large (small) x probes small (large) N. This is borne out by comparing F_2^{ep} and F_2^{en} (Fig. 7). At large x, the fact that some of the partons must be different in p and n is significant and $F_2^{ep} \neq F_2^{en}$. At small x, however, the parton multiplicity is high and the difference is insignificant.

4.3 $\nu N \to \mu + \ldots$

We now assume that partons are quarks and, for simplicity, consider just three flavours, u, d and s, and put $\theta_c = 0$. The basic processes are

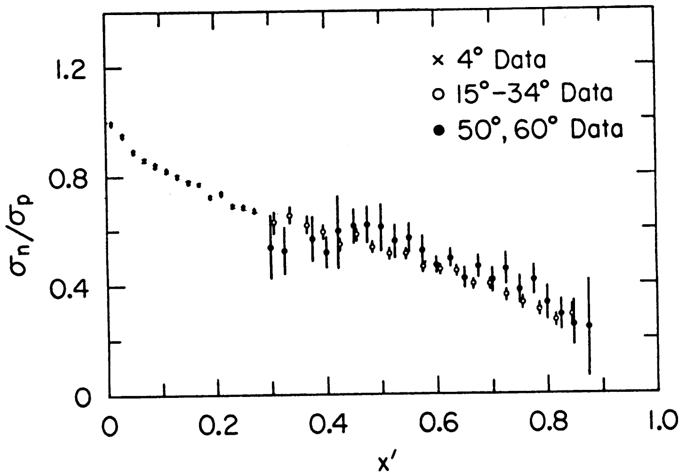


Fig.7. F_2^{en}/F_2^{ep} vs. x'. Data from SLAC-MIT experiments[36] (the Q^2 range of the data is roughly 2-20 GeV^2).

$$\nu d \to \mu^- u, \qquad \nu \bar{u} \to \mu^- \bar{d}$$

$$\bar{\nu} u \to \mu^+ d, \qquad \bar{\nu}\bar{d} \to \mu^+ \bar{u}$$

and the elementary cross-sections to be fed into eqn. 4.5 are

$$\frac{d\sigma^{\nu Q,\bar{\nu}\bar{Q}}}{dQ^2} = \frac{G^2}{\pi} \tag{4.7}$$
$$\frac{d\sigma^{\bar{\nu}Q,\nu\bar{Q}}}{dQ^2} = \frac{G^2}{\pi}(1-y)^2$$

Remarks:

a) Since Fermi's constant $G = \frac{10^{-5}}{M_p^2}$ these formulae have the right dimensions and in fact it is necessary that $\frac{d\sigma}{dQ^2} = G^2 g(y=Q^2/S_i)$.

b) The y dependence is easy to understand from helicity conservation. Consider $\bar{\nu}Q$ scattering in the $\bar{\nu}Q$ centre of mass frame where $2y=1-\cos\theta_{c.m.}$ (Fig.8). Initially $J_z=+1$ along the $\bar{\nu}$ direction; finally $J_z=+1$ along the μ^+ direction (assuming a V-A quark current and neglecting the quark and muon masses).

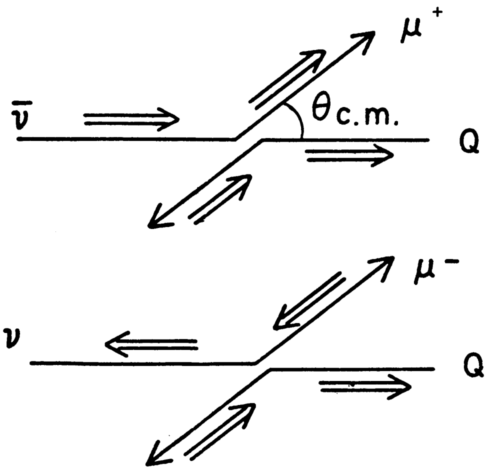


Fig.8. $\bar{\nu}Q$ and νQ scattering. The arrows represent the directions of the spins.

This process is therefore forbidden when θc.m.$=\pi$ or y=1 (explicit calculation is needed to find the power of 1-y). νQ scattering, however, is not forbidden at any angle and the process is isotropic.

c) With right-handed couplings to new quarks, the helicity factors for ν and $\bar{\nu}$ must obviously be exchanged and

$$\frac{d\sigma^{[\nu Q,\bar{\nu}\bar{Q}]_{RH}}}{dy} \sim (1-y)^2$$
$$\frac{d\sigma^{[\bar{\nu}Q,\nu\bar{Q}]_{RH}}}{dy} \sim 1$$

We denote the distribution of up, down, ... quarks in the proton by u(x), d(x) ... The distributions $u^n(x)$, $d^n(x)$ in the neutron are related by charge symmetry of the strong interactions

$$u^n(x) = d(x)$$
$$d^n(x) = u(x) \text{ etc.}$$

On isoscalar targets, neutrinos therefore measure

$$d(x) + d^n(x) = d(x) + u(x) \equiv Q(x)$$

and

$$\bar{u}(x) + \bar{u}^n(x) = \bar{u}(x) + \bar{d}(x) \equiv \bar{Q}(x)$$

and antineutrinos measure the same quantities. Combining eqns. 4.7 and 4.1 we find:

$$\frac{d^2\sigma^\nu}{dxdy} = \frac{G^2 s}{2\pi}\left[xQ(x) + x\bar{Q}(x)(1-y)^2\right]$$
$$\frac{d^2\sigma^{\bar{\nu}}}{dxdy} = \frac{G^2 s}{2\pi}\left[x\bar{Q}(x) + xQ(x)(1-y)^2\right] \tag{4.8}$$

A simple first guess would be to assume that the nucleon consists of just three quarks. This would imply

a) $\frac{d\sigma^\nu}{dy} = c$, $\frac{d\sigma^{\bar{\nu}}}{dy} = c(1-y)^2$

b) $\sigma^{\bar{\nu}}/\sigma^\nu = 1/3$ (equivalent to a), assuming scaling and charge symmetry of the weak current).

c) $\frac{\sigma^{\nu n}}{\sigma^{\nu p}} \simeq 2$ (assuming $\frac{u(x)}{d(x)} \simeq 2$).

d) $\sigma^\nu = G^2 s/2\pi \simeq 1.5 E_\nu \times 10^{-38} cm^2 GeV^{-1}$

since the momentum conservation sum rule

$$\sum_i \int dx\, x f_i(x) = 1$$

implies

$$\int dx\, x Q(x) = 1$$

if there are just three quarks in the nucleon.

The first three results are in approximate agreement with the data[37,38,39,30], although some $\bar{Q}$ contribution is required by the y distributions at small x and the concomitant fact that $\frac{\sigma^{\bar{\nu}}}{\sigma^\nu}$ is slightly bigger than 1/3 (the data give $\int x\bar{Q}/\int xQ \simeq 0.05-0.10$). However, the assumption that $\int xQ(x)dx=1$ overestimates σ^{ν} by a factor of about 2. This suggests that about one half of the proton's momentum is carried by gluons, which do not interact weakly or electromagnetically, which are needed to mediate the strong interquark force in any renormalizable theory (they were first introduced[40]) in the parton model to explain the fact that F_2^e is implausibly small even if the partons are non-integrally charged quarks).

One experiment[41]) claimed very substantial energy dependent deviations from this simple picture but this is not supported by recent data from experiments[38,30]) with much higher statistics and bigger acceptance (couplings to new quarks were invoked to explain these deviations; we discuss their effects in Section 7). However, some deviations are certainly expected. First, the effects of $\Delta C=1$ currents should be included and a non-zero Cabibbo angle should be used; this is considered in Section 5. Second, scaling violations will produce changes; this is the subject of Section 6.

4.4 General Tests of the Quark Parton Model

We now turn to some more general and incisive tests of the quark parton model, which have the virtue that they survive the introduction of the scaling violations predicted by asymptotically free gauge theories (see Section 6).

First it is necessary to obtain the parton model formulae for the σ_i or F_i. Comparison of the parton formulae (4.8) with the general form of the differential cross-section given in Section 2 gives

$$\sigma_0 = 0 \quad \text{or} \quad 2xF_1 = F_2$$
$$F_2^{\nu p} = 2x\,[d(x) + \bar{u}(x)]$$
$$F_2^{\nu n} = 2x\,[u(x) + \bar{d}(x)]$$
$$F_3^{\nu p} = 2x\,[\bar{u}(x) - d(x)]$$
$$F_3^{\nu n} = 2x\,[\bar{d}(x) - u(x)]$$

(keeping to three flavours of quarks and $\theta_c=0$ for simplicity; it is a simple exercise to write the results in the general case). The fact that $\sigma^{\bar{\nu}}/\sigma^{\nu}$ is close to one third requires that $\sigma_0^{\nu} \ll \sigma_L^{\nu}$ or $2 \times F_1 \simeq F_2$ (assuming scaling and charge symmetry). Combined with eqn.4.6, the parton model formulae lead to a number of important results for the remaining structure functions F_2 and F_3:

a) Since there are six measurable structure functions ($F_2^{\nu p,\nu n}$, $F_2^{ep,en}$, $F_3^{\nu p,\nu n}$) which are determined in terms of five quark distributions (u, d, $\bar{u}$, $\bar{d}$, $s+\bar{s}$) there is one relation[40])

$$6\,(F_2^{ep} - F_2^{en}) = x\,(F_3^{\nu p} - F_3^{\nu n}).$$

Unfortunately this is very hard to test.

b) The model requires[40])

$$F_2^{\nu(p+n)} \leqslant \frac{18}{5}\,F_2^{e(p+n)}$$

since the right-hand side differs from the left by $\frac{2}{5}(s+\bar{s})$ which is positive semi-definite. The Gargamelle data, displayed in Fig.9, show that the inequality is approximately saturated for $x>0.1$ (where we would expect to find few partons, and hence $s+\bar{s}$ to be small, according to the arguments above). In passing we remark that F_3 can be extracted from the Gargamelle data as well as F_2 and they can be combined to extract the functions Q and $\bar{Q}$

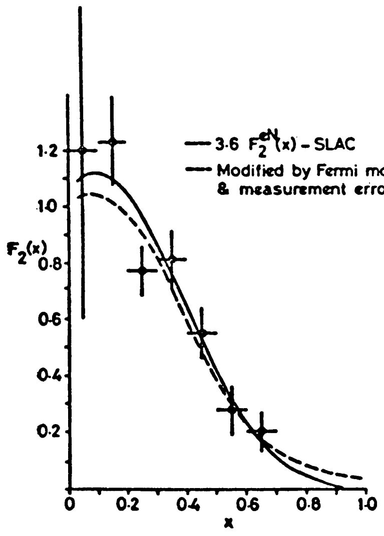


Fig.9. $F_2^{\nu N}$ vs. x measured in the Gargamelle experiment[37]) for $Q^2 \gtrsim 1$ compared to the SLAC-MIT data for 18/5 F_2^{eN} in the same Q^2 range.

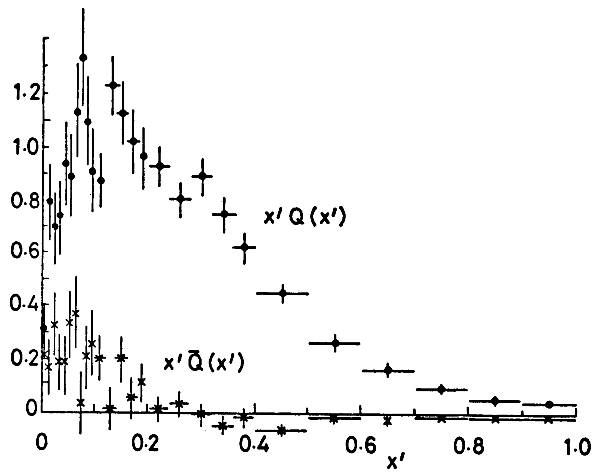


Fig.10. Quark and anti-quark momentum distributions as a function of x' as measured in the Gargamelle experiment[37]).

shown in Fig.10. They show exactly the behaviour anticipated although Q^2 is very small at small x so this should not necessarily be taken too seriously.

c) The model leads to other interesting inequalities such as[42])

$$4 \geqslant \frac{F_2^{en}}{F_2^{ep}} = \frac{4(d+\bar{d}) + u + \bar{u} + s + \bar{s}}{4(u+\bar{u}) + d + \bar{d} + s + \bar{s}} \geqslant \frac{1}{4}$$

(A complete set follows from the 'isospin inequalities' $2u(x)>d(x)$, $2\bar{d}(x)>\bar{u}(x)$ and $q,\bar{q}>0$).

d) The parton distributions must be consistent with the quantum numbers (B, Y, I_3, S, C ...) of the target. For example, there must be three more quarks than antiquarks in the proton, i.e.

$$3 = \int dx\left[u+d+s+c-\bar{u}-\bar{d}-\bar{s}-\bar{c}\right].$$

Using the sum rules

$$S = \int dx(s-\bar{s}) = 0$$

$$C = \int dx(c-\bar{c}) = 0, \text{ etc.}$$

we can express the right-hand side in terms of F_3 thus[43)]

$$-\int F_3^{\nu(p+n)}\,dx = 6$$

The data are shown in Fig.11. Unforunately the small x contribution (which is important) corresponds to small Q^2 so the agreement should be taken too seriously but tests at much higher Q^2 will shortly become available. The other sum rules which can be derived in this way (which are catalogued in ref.1) are much harder to test since they involve the difference of structure functions on neutrons and protons.

e) Finally we note that energy momentum conservation requires[1,44)]

$$1 = \int dx\; x(u+d+s+\bar{u}+\bar{d}+\bar{s}) + \varepsilon$$

where ε is the fraction of the proton's momentum carried by gluons or other flavours of quarks. Still neglecting charm and keeping $\theta_c=0$ we can write

$$\varepsilon = 1 + \int dx\left(\frac{3}{4}F_2^{\nu(p+n)} - \frac{9}{5}F_2^{e(p+n)}\right)$$

Experimentally $\varepsilon \simeq 0.5$.

To summarise: the data support the quark-parton model although more stringent tests of the sum rules are desirable. The fact that $F_2{}^\nu/F_2{}^e$ is given correctly by the quark model at large x strongly supports the picture of incoherent scattering off (effectively) non-integrally charged particles. Since non-integral quark charges give the correct ratio we would obviously expect to obtain the wrong result with integral charges. Equivalently we can say that the data favour non-integral hypercharge. Using CVC we can write

$$\frac{F_2^{eN\to I=0}}{F_2^{eN\to I=1}} = \frac{\langle Y^2\rangle_{eff.}}{4\langle I_3^2\rangle_{eff.}} = \frac{4F_2^{eN}-F_2^{\nu N}}{F_2^{\nu N}}$$

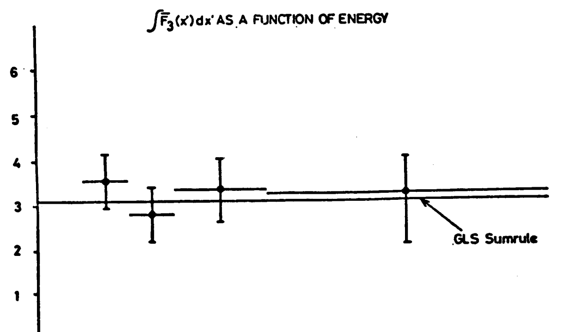


Fig.11. $-\frac{1}{2}\int F_3^{\nu n+\nu p}dx'$ from the Gargamelle experiment[37)] as a function of the energy ($\langle Q^2\rangle \sim 0.2E$).

The data show that this ratio is small. It can be argued that this is exactly as expected, given continuity between photo-production and large Q^2 electroproduction, since the fairly successful, ρ, ω, ϕ dominance model of photo-production gives

$$\frac{\sigma(\gamma N\to I=0)}{\sigma(\gamma N\to I=1)} \simeq 1/9$$

However, I do not think that this belittles the success of the quark-parton model. Rather it shows that photo-production already suggested non-integral hypercharge for underlying constituents.

5. Charmed Particle Production and Di-Leptons

The GIM current (Section 1.1) leads to the following predictions in the framework of the quark-parton model for an I=0 target

$$\left.\frac{d\sigma^\nu}{dx\,dy}\right|_{\Delta C=1} = \frac{G^2 s}{2\pi}\left[x(d+u)\sin^2\theta_c + 2xs\cos^2\theta_c\right]$$

$$\left.\frac{d\sigma^{\bar\nu}}{dx\,dy}\right|_{\Delta C=1} = \frac{G^2 s}{2\pi}\left[x(\bar{d}+\bar{u})\sin^2\theta_c + 2x\bar{s}\cos^2\theta_c\right]$$

Charged current data give

$$\frac{\int x\bar{Q}\,dx}{\int xQ\,dx} \simeq 0.05 - 0.10$$

and it seems reasonable to suppose that

$$\frac{2\int xs\,dx}{\int xQ\,dx} = \frac{2\int x\bar{s}\,dx}{\int xQ\,dx} \lesssim \frac{\int x\bar{Q}\,dx}{\int xQ\,dx}$$

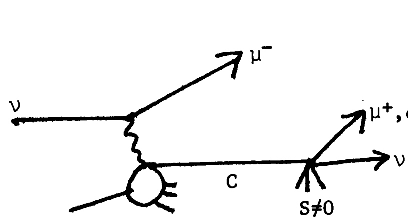
Since $\tan^2\theta_c$ is about 0.06, we see that the two contributions to $d\sigma^\nu$ are expected to be of the same order of magnitude. In fact the model makes the following definite predictions for $d\sigma^{\nu,\bar\nu}_{\Delta C=1}$

a) the y distribution should be flat for ν and $\bar\nu$.

b) $d\sigma^\nu_{\Delta C=1}$ should have two components, one with the same x distribution as the $\Delta C=0$ cross-section contributing about 6% (relative to $\sigma^\nu_{\Delta C=0}$) and the other, concentrated at small x, contributing less than or of order 5 to 10%.

c) $d\sigma^{\bar\nu}_{\Delta C=1}$ should be concentrated at small x and is expected to be less than or of order 5 to 10% relative to $\sigma^\nu_{\Delta C=0}$ or less than or of order 12 to 25% relative to $\sigma^{\bar\nu}_{\Delta C=0}$.

It might be possible to detect these components with their characteristic threshold behaviour in the inclusive cross-sections (especially the additional flat contribution to $\frac{d\sigma^{\bar\nu}}{dy}$ at small x), but they are small and may be masked by scaling violations. A cleaner signature is provided by the semileptonic decays of the charmed particle



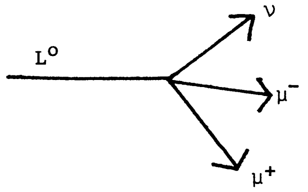
We know from e^+e^- annihilation experiments that charmed particles decay to μ or e 10-20% of the time each[45]. We therefore anticipate

$$\frac{\nu \to \mu^-\mu^+}{\nu \to \mu^-} \simeq (1-2)\% \quad \text{and} \quad \frac{\bar{\nu} \to \mu^+\mu^-}{\bar{\nu} \to \mu^+} \approx (2-3)\%$$

(in the absence of any experimental cuts). In fact $\nu \to \mu\bar{\mu}^+$ ... and $\nu \to \mu e$ ... events have been observed [46,47] with the following features

a) $$\frac{\langle E_{\mu^-}\rangle^{\nu}}{\langle E_{\mu^+}\rangle^{\nu}} \quad \text{and} \quad \frac{\langle E_{\mu^+}\rangle^{\bar{\nu}}}{\langle E_{\mu^-}\rangle^{\bar{\nu}}}$$

are considerably greater than one (because there is an experimental muon energy cut it is not sensible to quote a value without giving details of each experiment). This rules out[48,49] the notion that the muons might both be the decay products of a single neutral heavy lepton



which would require[49]

$$2 > \langle E_{\mu^-}\rangle / \langle E_{\mu^+}\rangle > \tfrac{1}{2}$$

for any mixture of V and A currents (for any mixture of S, P, T, V and A currents $2 \to 2.09$). However, models in which the second muon is the decay product of a heavy hadron account for the energy spectrum very well[50].

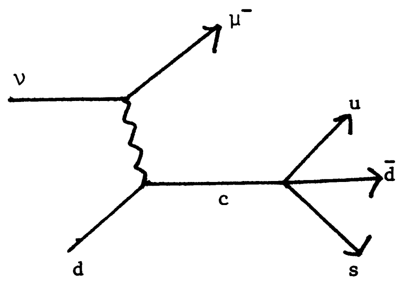
b) There is a strong anticorrelation between the two muons in azimuthal angle which increases with the energy of the second muon[47]. This is a very strong indication that the second muon is associated with the hadron shower. It is well reproduced by models in which the second muon comes from the decay of a heavy hadron.

c) The x and y distributions of the $\mu^-(\mu^+)$ in $\nu(\bar{\nu})$ induced dimuon events is exactly that anticipated by the charm model[47] discussed above. Furthermore, the experimental cuts can be applied to the model in which case it can successfully account for the observed rate (the cuts have a big effect since the expected spectrum of the second muon is soft - especially for $\sigma_{\bar{\nu} \to \mu\mu}$).

d) $\nu \to \mu^- e^+$ events have been observed[51] in bubble chamber experiments at a rate consistent with that for $\nu \to \mu^-\mu^+$. Very strong evidence that charmed particle production is responsible is provided by the observation of a strong correlation between these events and K^0_S or Λ production. There is considerable fluctuation between the strength of the correlation observed in different experiments (which still have quite small statistics) but it seems to be compatible with the production of about one strange particle per $\mu^- e^+$ event as expected if the GIM current is responsible.

e) $\nu + Fe \to \mu^-\mu^-$ ... and $\bar{\nu} + Fe \to \mu^+\mu^+$ ... events have been observed, most of which can be attributed to π or K decay in flight[52]. After subtracting the best estimate of the background, a signal of $(3 \pm 2) \times 10^{-4} \sigma_{\nu \to \mu, \bar{\nu} \to \bar{\mu}}$ remains. If we take the non-zero value serious, it is consistent with the production of charmed-anticharmed particle pairs in about 0.5-1% of the charged current events, which does not seem unreasonable.

Further evidence for charmed particle production by neutrinos is provided by the observation of a $\Delta S = -\Delta Q$ event $\nu p \to \mu^- \Lambda^0 \pi^+ \pi^+ \pi^+ \pi^-$ in a hydrogen bubble chamber experiment at Brookhaven[53]. In terms of quarks, apparent $\Delta S = -\Delta Q$ events are produced by the sequence



The observed event can be interpreted as the production of a charmed baryon (calculations of the cross-section for charmed baryon production may be found in ref.20).

In addition an event has been observed in a neutrino emulsion experiment at FNAL[18] in which a V^0 is produced and one charged particle travels 182μm before decaying, corresponding to a lifetime of about 10^{-13}secs as expected for a charmed particle.

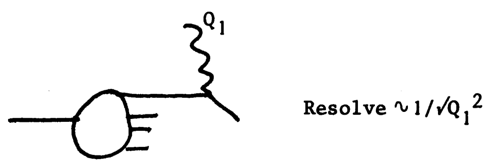
6. Scaling Violations

6.1 Preliminary Remark

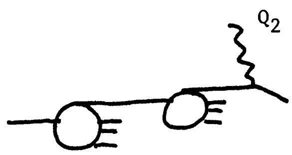
It is important to realize that the average Q^2 in neutrino experiments at FNAL and the CERN SPS is about the same as the maximum Q^2 for which the $\mu N \to \mu$ ... data are available. Therefore we cannot use the scaling violations observed in the electromagnetic structure functions to predict the effects in neutrino experiments without making considerable extrapolations, for which purpose a model is needed. I will use an asymptotically free gauge theory of the strong interactions[54], which is easily the most attractive model available. However, rather than follow the formal approach, which requires a considerable knowledge of renormalization theory, I will discuss a simple physical interpretation of the results developed (before some of the formal approach!) by Kogut and Susskind[55,56], following ideas due to Kadanoff, Wilson and Polyakov.

6.2 The Scale Invariant Parton Model

When virtual photons of momentum Q_1 are scattered from a target they are able to resolve constituents of size proportional to $1/\sqrt{Q_1}^2$ (for fixed $x \neq 0$) but are not sensitive to substructure on a much smaller scale:

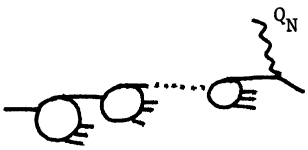


When the momentum is increased, however, substructure may be revealed



Resolve $\sim 1/\sqrt{Q_2}^{\,2}$

and so on



Resolve $\sim 1/\sqrt{Q_N}^{\,2}$

This process of successively resolving clusters made of smaller clusters which are made of smaller clusters, etc. is easiest to visualize when there is a relatively sharp transition between length scales, as occurs in the successive resolution of molecules, atoms, nuclei, nucleons, quarks .... However, quark structure, consisting of virtual gluons and Q-$\bar{Q}$ pairs, is unveiled continuously in a quark field theory as quarks are probed at shorter and shorter distances.

In order to discuss scattering at momentum Q_N we need an N-dependent effective Hamiltonian H_N which describes the $(N-1)^{th}$ layer in terms of the N^{th} layer constituents, e.g. at the stage at which quasi-elastic scattering from single nucleons dominates we need a Hamiltonian which describes the nucleus in terms of nucleons but the existence of nucleon substructure is irrelevant. At the quark level, the nature of the constituents involved does not vary with Q^2 but the Hamiltonian $H(Q^2)$ will depend on an effective coupling constant $g(Q^2)$ which changes with Q^2 as different sizes of cluster are resolved. We consider three possibilities as $Q^2 \to \infty$:

a) $g^2 \to \frac{1}{(Q^2)^P}$. In this case the constituents would behave like free particles at large Q^2 and the naive parton model would be valid. This can only occur in superrenormalizable theories in which there is a scale in the interaction (e.g. in a scalar field theory with $H_{int} = \lambda\phi^3$ where λ has the dimensions of a mass), which sets the size scale below which interactions can be neglected. However, there are no superrenormalizable theories involving spin 1/2 particles. In renormalizable theories like QED or QCD the coupling constant is dimensionless which means that there is no distance scale below which interactions can be neglected; a virtual photon never has enough resolving power to reveal totally naked partons. However, by imposing a transverse momentum cut off on a renormalizable theory (which spoils the locality and Lorentz covariance), it can be made to behave like a superrenormalizable theory and the parton picture will be valid for $Q^2 \gg (\vec{p}_T^{\,2})$ max; this was the basis of an interesting investigation of parton models due to Drell, Levy and Yan[57].

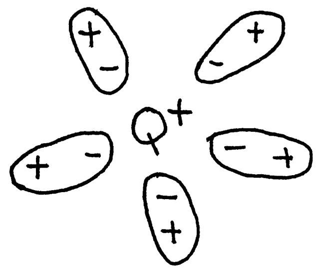
b) $g^2 \to g_*^2 \neq 0$. This "fixed-point" behaviour could occur in any field theory. If it occurs, the relation of one layer to the next becomes universal in a sense we shall discuss below.

c) $g^2 \to 1/\ln Q^2$. This is asymptotic freedom. It can only occur in non-Abelian gauge theories[54].

Before exploring the consequences of the Q^2 dependent parton picture corresponding to cases b) and c), we consider the force between two test charges in a di-electric which illustrates the concept of an effective distance dependent coupling constant. The potential is

$$V = \frac{Q_1 Q_2}{4\pi\epsilon r}$$

ε differs from ε_o due to polarization



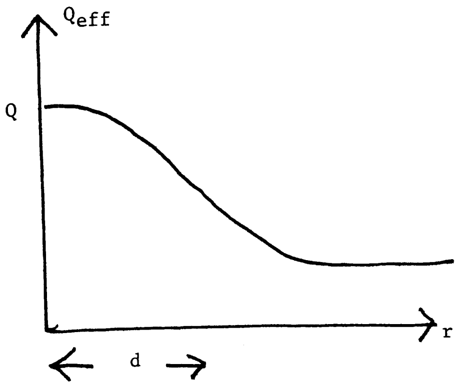
which shields the charge. For r very much less than the molecular spacing d, however, there will be no shielding and $\varepsilon \to \varepsilon_o$. We can define an effective charge

$$Q_{eff.}(r) = Q\sqrt{\frac{\epsilon_o}{\epsilon}}$$

in terms of which

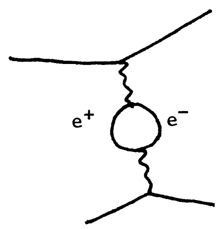
$$V = \frac{Q^1_{eff.}\, Q^2_{eff}}{4\pi\epsilon_o r}$$

which behaves as

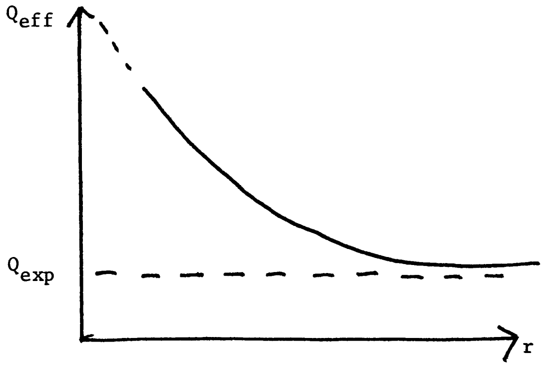


(Note that in a metal there is infinite shielding and $Q_{eff} \to 0$ as $r \to \infty$).

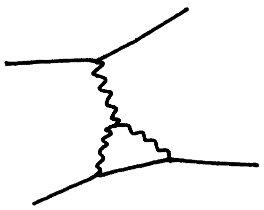
In quantum electrodynamics vacuum polarization diagrams like



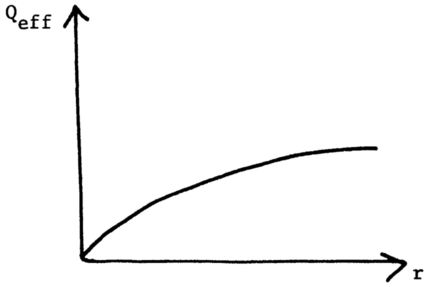
give rise to shielding so that



In non-Abelian gauge theories, however, there are additional diagrams such as

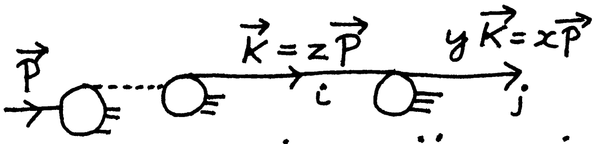


among others. Detailed calculation shows that there is antishielding ($\varepsilon<\varepsilon_0$) in this case. In fact the antishielding is infinite at infinite distances so we need $Q_{eff}(0)=0$



No-one has yet succeeded in "explaining" this result physically (we cannot even say which diagrams are responsible, since their relative contributions are gauge dependent, but we might imagine that the diagram above "spreads out" the charge density making it vanish as $r\to 0$).

Keeping a discrete index N for simplicity, we define probability distributions for the partons in the following way



Prob $\quad F^i_{N-1}(z),\ f^{ij}(y),\ F^j_N(x)$

In terms of these distributions, this picture has the following implications:

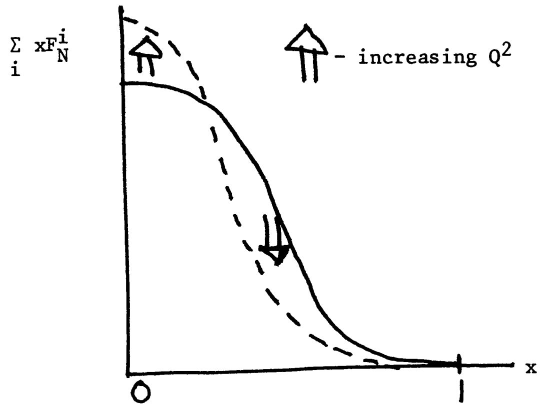
a) Momentum conservation reads

$$\sum_i \int x\, F^i_N(x)\,dx = 1\,.$$

However, since the number of partons increases indefinitely with N we expect

$$\langle x\rangle_N = \sum_i \int x^2 F^i_N(x)\,dx \xrightarrow[N\to\infty]{} 0$$

Thus $\Sigma x F^i_N$ must behave as



Since $F_2{}^e=\sum_i Q_i^2 x F_N{}^i$, the picture predicts that $F_2{}^e$ also increases at small x and decreases at large x with increasing Q^2 - a trend seen in the data.

b) Since the distributions must yield the correct conserved quantum numbers B, I_3, Q ... for all N, the parton sum rules and inequalities will be satisfied.

c) The probabilities F_{N-1}, F_N and f_N are obviously related by

$$F_N(x) = \int_0^1 F_{N-1}(x/y)\, f_N(y)\,\frac{dy}{y}$$

Taking moments of this equation we obtain

$$\frac{\int x^p F_N(x)\,dx}{\int x^p F_{N-1}(x)\,dx} = \int y^p f_N(y)\,dy$$

Since specific models will tell us something about f, we see that it is the behaviour of the moments of the structure function which will be predicted directly.

In the case of a fixed point theory ($g_N^2 \to g_*^2$)

$$\int y^p f_N(y)\,dy = q_p$$

is independent of N (note that $q_1=1$, $q_p<1$ for $p>1$, if there is a single species of constituent). The ratio of the size of successive clusters is also N-independent

$$\frac{R_N}{R_{N+1}} = \Lambda$$

and we have

$$\frac{R_0}{R_N} = \sqrt{\frac{Q^2}{Q_0^2}} = \Lambda^N$$

N is therefore related to Q^2 by

$$N = \frac{1}{2}\,\frac{\ln Q^2/Q_0^2}{\ln \Lambda}$$

and

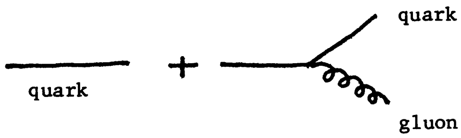
$$\frac{\int x^p F_N(x)dx}{\int x^p F_0(x)dx} = (q_p)^N = \left(\frac{Q_0^2}{Q^2}\right)^{\frac{\ln q_p}{2\ln\Lambda}}$$

i.e. the moments decrease as powers of Q^2 in a fixed point theory.

In the case of non-Abelian gauge theories which are asymptotically free

$$g^2(t) = c/t$$

where $t=\ln Q^2/\mu^2$ and c is calculcable. Since g is small for large Q^2, the function $f(y,t)$ which relates the distribution F at t and $t+\delta t$ can be calculated perturbatively from graphs like[58]



Hence $f(y,t)=\delta(1-y)+g^2(t)k(y)$ where k is known. The F's therefore satisfy

$$F(x,t+\delta t) = F(x,t) + g^2(t)\int F(x/y,t)k(y)\frac{dy}{y}\delta t$$

or

$$\frac{dF(x,t)}{dt} = g^2(t)\int F(x/y,t)k(y)\frac{dy}{y}$$

Taking moments

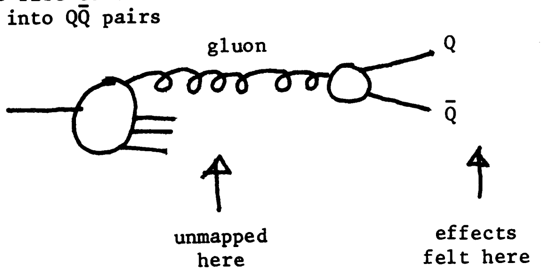
$$\frac{d}{dt}\int x^p F(x,t)dx = \frac{A_p}{t}\int x^p F(x,t)dx$$

where $A_p = c\int y^p k(y)dy$ is known.
The solution of this differential equation is

$$\frac{\int x^p F(x,t)dx}{\int x^p F(x,t_0)dx} = \left(\frac{t}{t_0}\right)^{A_p}$$

i.e. the moments of the structure functions decrease as calculable powers of log Q^2 in asymptotically free gauge theories.

We have ignored a complication due to the fact that the function f is actually a matrix f_{ij} for the probability of finding parton type j in parton type i. The most annoying consequence of this coupling between different species is that the neutral gluons, whose effects are not felt at t, give rise to effects at $t+\delta t$ since they can dissociate into $Q\bar{Q}$ pairs



The effects of gluons cancel out in the case of the so-called non-singlet structure functions such as $F_2^{ep}-F_2^{en}$ (since the glue is isoscalar) and F_3 (which is proportional to $Q(x,t)-\bar{Q}(x,t)$). In general, however, it means that data are needed at two different values of t as boundary conditions to solve the differential equations. The solution of this mixing problem reveals not only that truly asymptotically all species of quark and antiquark are equally probable (as expected infinitely far down the quark fragmentation chain) but that the momentum is shared between quarks/antiquarks and gluons in the ratio 3m/16 in a model with m flavours and three colours. This gives rise to sum rules[54] for the limits of $\int F_2^{e,\nu}dx$, as $t\to\infty$. It turns out that these limits are not expected to be reached until astronomical values of Q^2, but it is important to note that the asymptotic limits are below the measured values of these moments which are therefore expected to decrease with Q^2.

6.3 Predictions of Asymptotically Free Gauge Theories

We have seen that the parton picture can be modified to accommodate asymptotic freedom (so that its successes survive) by letting the distributions change with log Q^2 in a well defined way. The predicted Q^2 variation is compatible with the scaling violations observed in $\mu(e)N\to\mu(e)$ ... (with the unknown scale parameter μ^2 probably being in the range 0.05 to 0.25 GeV^2), but it is not yet definitely tested[59].

Turning to σ^ν, let us use eqn.4.9, neglecting $\bar{Q}$ in a first qualitative discussion, so that

$$\frac{\sigma^\nu}{E} \sim \int_0^1 dx \int_0^1 dy\, x\,Q(x, Q^2 = 2xyME_\nu)$$

Note that at small x (where Q rises with Q^2), Q^2 is small but at large x (where Q falls) Q^2 is large. It follows that the change in shape of $Q(x,Q^2)$ as a function of Q^2 would cause $\frac{\sigma^\nu}{E}$ to fall with E_ν even if

$$\langle xQ(x,Q^2)\rangle \equiv \int xQ(x,Q^2)\,dx$$

were independent of Q^2. In fact it follows from the momentum sum rule of asymptotic freedom that $\langle xQ\rangle$ decreases towards a constant value (reached at astronomical Q^2) and $\langle x\bar{Q}(Q^2)\rangle$ rises to the same value. This increases the tendency of σ^ν/E to fall[60].

Any discussion of $\sigma^{\bar{\nu}}$ is a more delicate matter requiring detailed calculations since $\bar{Q}$ plays a more important role and the contribution of Q is suppressed at large Q^2 by the $(1-y)^2$ factor. In fact it turns out that asymptotic freedom tends to make $\frac{\sigma^{\bar{\nu}}}{E}$ fall also[60].

The results of a model calculation[60] are shown in Figs.12 and 13. Remarks:

a) If we separate the $\Delta C=0$ contribution we get

E	50	150	400
$\left(\frac{\sigma^\nu}{E}\right)_{\Delta C=0}$	0.52	0.48	0.46
$\left(\frac{\sigma^{\bar{\nu}}}{E}\right)_{\Delta C=0}$	0.24	0.22	0.22

This not only allows us to estimate the importance of charmed particle production but also to isolate the effects of asymptotic freedom which are seen to

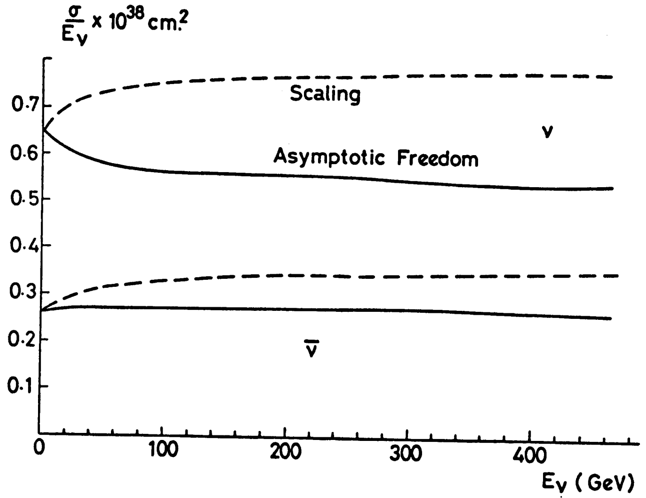


Fig. 12. $\frac{\sigma^{\nu,\bar{\nu}}}{E}$ vs. E assuming scaling (----) (the rise is due to charmed particle production) and asymptotic freedom (——) (ref.60).

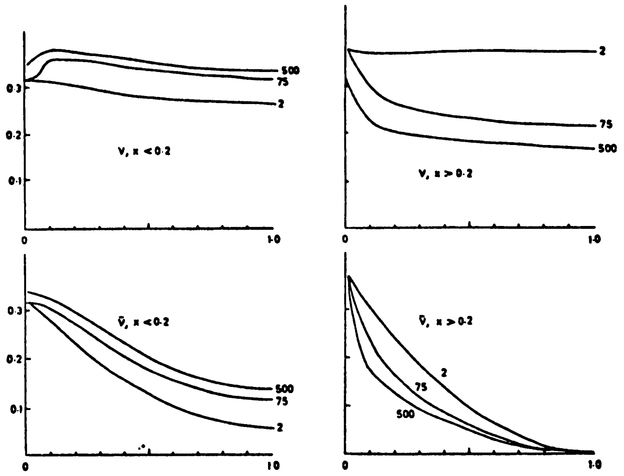


Fig. 13. $\frac{1}{E}\frac{d\sigma^{\nu,\bar{\nu}}}{dy}$ vs. y for x>0.2 and x<0.2 according to asymptotic freedom[60].

make both $\frac{\sigma^\nu}{E}$ and $\frac{\sigma^{\bar{\nu}}}{E}$ fall with energy. It also turns out that for x>0.2 charmed particle production contributes negligibly to $d\sigma^{\bar{\nu}}$ and gives about 5% ($\tan^2\theta_c$) of $d\sigma^\nu$ as expected; for x<0.2, where Q^2 is relatively small, almost all the energy dependence is due to charmed particle production.

b) $\sigma^{\nu,\bar{\nu}}/E$ vary by less than 5% between 50 and 500 GeV according to asymptotic freedom (unless there are more than four quarks or M_W is unexpectedly small; the possibility of conspiracies between such effects is analysed in ref.60).

c) Note the manifest scaling violation in $\frac{d\sigma^{\bar{\nu}}}{dy}$ for x>0.2 which does not have the form $a+b(1-y)^2$.

The first measurements of σ^ν/E^ν from the CERN SPS do show a fall between 10 and 100 GeV[30] and it is consistent with the predictions of asymptotic freedom. This is very encouraging. However, much more data and work will be needed to show whether it is due to asymptotic freedom. Measurement of the total cross-section is a crude way to study scaling violations. But if scaling is violated, the old analyses in terms of x and y distributions are not quite so useful and it is really necessary to try to separate the structure functions as functions of x and Q^2 which is a difficult job.

7. New Quarks

Reasons why new quarks might exist include

a) Why not? (In fact the Υ particle[61], discovered while these lectures were being given, is probably evidence for a new quark[62]).

b) A desire for symmetry between the leptons, of which there are now probably six doublets with left-handed weak couplings

$$\begin{pmatrix}\nu_e\\ e\end{pmatrix}_L \qquad \begin{pmatrix}\nu_\mu\\ \mu\end{pmatrix}_L \qquad \begin{pmatrix}\nu_\tau\\ \tau\end{pmatrix}_L$$

and the quarks, which might be organized into doublets with left-handed couplings to the weak current, thus

$$\begin{pmatrix}u\\ d_c=d\cos\theta_c+s\sin\theta_c\end{pmatrix}_L , \begin{pmatrix}c\\ s_c=s\cos\theta_c-d\sin\theta_c\end{pmatrix}_L , \begin{pmatrix}t\\ b\end{pmatrix}_L$$

(given the existence of τ and ν_τ, some further quarks or leptons are needed to cancel the triangle anomalies which spoil renormalizability in gauge theories; this arrangement does this successfully).

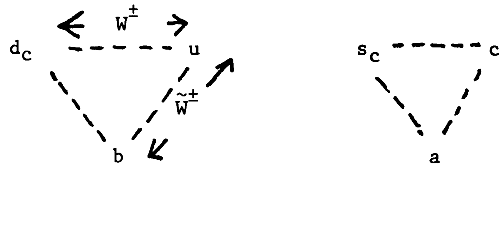
c) A desire for every quark to enjoy left and right handed couplings in equal measure, so that the underlying Lagrangian can be parity conserving. Since u↔d and u↔s transitions are known to be predominantly left handed, this requires new quarks to make up right handed doublets e.g. with two new quarks

$$\begin{pmatrix}u\\ b\end{pmatrix}_R , \begin{pmatrix}t\\ d\end{pmatrix}_R , \begin{pmatrix}c\\ s\end{pmatrix}_R .$$

d) The representations of the SU(2)xU(1) weak gauge group might be larger than doublets which would require quarks of charge +5/2 and/or −4/3 in weak multiplets such as

$$\begin{pmatrix}X\\ u\\ d_c\\ Y\end{pmatrix}$$

e) The gauge group SU(2)xU(1) might be embedded in a bigger group such as SU(3)xU(1) which would require additional quarks e.g. with quarks in left handed triplets we might have



(in this case the new quarks would be coupled to the old via new currents and would be produced by neutrinos in association with new leptons, in the absence of substantial mixing).

We will concentrate on new quark production off valence quarks which could make substantial contributions to $\sigma^{\nu,\bar{\nu}}$. We consider the optimal case in which the weak current acts with its full strength although the effects could, of course, be diluted by mixing angles[63]. In a naive three quark plus glue parton model the possibilities are:

Transition	Nature of Coupling	Contribution to $\frac{d\sigma}{dy}$	Change in $\sigma^{\nu N}$
$\nu d \to \mu^- q^{2/3}$	L(V-A)	Flat	$\sigma \to 2\sigma$
$\nu u \to \mu^- q^{5/3}$	R(V+A)	$(1-y)^2$	$\sigma \to \frac{4\sigma}{3}$
$\nu u \to \mu^- q^{-1/3}$	L(V-A)	$(1-y)^2$	$\sigma \to 2\sigma$
$\nu d \to \mu^- q^{-4/3}$	R(V+A)	Flat	$\sigma \to 4\sigma$

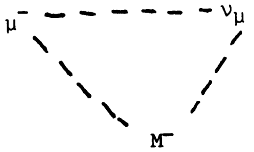
The degeneracy of the effects of $u \to q^{5/3}$ with $d \to q^{2/3}$ and $u \to q^{-1/3}$ with $d \to q^{-4/3}$ transitions in $\sigma^{\nu N}$ is of course lifted in $\sigma^{\nu p}$ and $\sigma^{\nu n}$. In all cases the effects are big enough to be readily detectable well above threshold and for this conclusion to survive the relaxation of the valence quark approximation and the introduction of the expected scaling violations. As well as changing $d\sigma^{\nu,\bar{\nu}}$, new quark production will contribute to $\nu/\bar{\nu} \to \mu^-\mu^+ \ldots$ assuming a non-negligible semi-leptonic branching ratio. The fact that $\frac{\sigma^{\bar{\nu}}}{\sigma^{\nu}}$ and $\frac{\sigma^{\nu \to \mu^-\mu^+}}{\sigma^{\bar{\nu} \to \mu^+\mu^-}}$ are nearly energy independent from 50 to 200 GeV allows us to conclude that if there are new quarks coupled to u and d with the full strength of the weak current their masses must be at least 10 GeV (the exact limit depends on the nature of the coupling and on how σ behaves near the threshold). Neutrino experiments are clearly a powerful diagnostic for new quarks and their weak couplings.

8. New Leptons

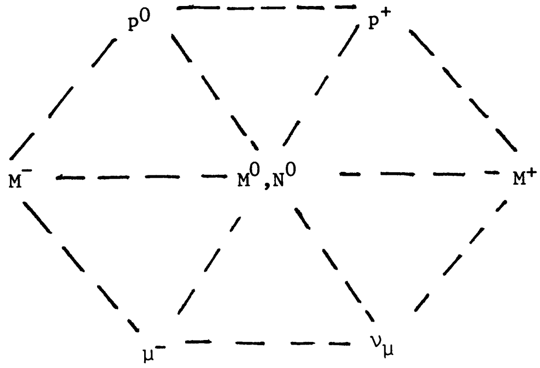
Most of the discussion above of why new quarks might exist applies also to new leptons[64]. As examplies of possible multiplet structures we could have SU(2)xU(1) triplets

$$\begin{pmatrix} M^+ \\ \nu_\mu \\ \mu^- \end{pmatrix}$$

allowing $\nu_\mu \to M^+$ transitions or SU(3)xU(1) triplets

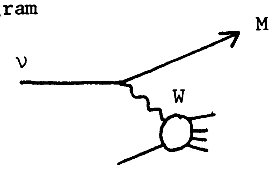


allowing $\nu_\mu \to M^-$ (presumably in association with new quark production, in the absence of mixing) or octets



allowing $\nu \to M^+$, M^0, N^0.

The production cross-section corresponding to the diagram



is easy to estimate assuming exact scaling and that the usual weak current is involved. The result is shown in Fig. 14.

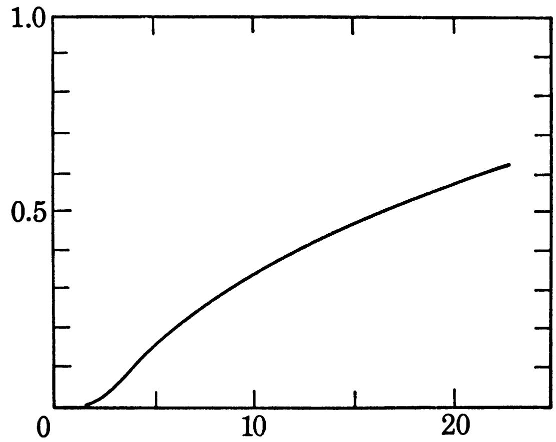


Fig. 14. $\frac{\sigma(\nu \to M)}{\sigma(\nu \to \mu)}$ vs. M^2/S assuming $g_{\nu M} = g_{\nu\mu}$ and exact scaling[65].

Possible final states and signatures for charged heavy lepton production are

Process	Signature
$\nu_\mu A \to M^+ \ldots$	
$\to \mu^+\nu\nu$	Apparent lepton non-conservation ($\to$limit[66]) $M_{M^+} > 8.4$ GeV if $g_{\nu M} = g_{\nu\mu}$)
$\to \nu$+hadrons	Apparent neutral current
$\to M^0 + \ldots$	

Process	Signature	
$\nu_\mu A \to M^- \ldots$		
$\to \mu^-\nu\nu$	Apparent small x, small y anomaly	
$\to e^-\nu\nu$	Anomalous e^- prod.	(rules out τ (1.9)$=M^-$ unless $g_{\nu M} << g_{\nu\mu}$)
$\to \nu$+hadrons	Apparent neutral current	
$\to M^0\mu^-\nu$		
$\hookrightarrow \mu^+\mu^-\nu$	Trilepton events	

A detailed discussion of branching ratios may be found, e.g., in ref.64. For present purposes we only need to note that leptonic modes allowed by selection rules are expected to have substantial branching ratios (of order 20%). Because of the spectacular leptonic signatures, current neutrino experiments are sensitive to leptons which couple to ν_μ with mass up to about 15 GeV unless $g_{\nu M} << g_{\nu\mu}$.

Recently some evidence for trilepton production $\nu_\mu \to \mu^-\mu^+\mu^- \ldots$ has been reported[67] which has produced a plethora of theoretical papers[68] interpreting this as evidence for the production of M^- and M^0 in the chain $\nu A \to M^- + \ldots$, $M^- \to M^0\mu^-\nu$, $M^0 \to \mu^+\mu^-\nu$ with $M_{M^-} \simeq 7$ GeV and $M_{M^0} \simeq 4$ GeV (and also as associated production of charmed particles[69], although this is harder to reconcile with some of the reported features of the data, or of joint new lepton and new quark production[70]). The CDHS experiment at the CERN SPS[71] has not confirmed the rate reported by the FNAL experiments, and although they have a few events it is even possible that they are due to background. However, the rate originally reported required $g_{\nu M^-} \simeq g_{\nu\mu^-}$ (with plausible assumptions for branching ratios) which would have had implications in gauge theories which are sufficiently interesting that they are worth discussing briefly.

Whereas $g_{\nu_\mu M^+} = g_{\nu_\mu \mu^-}$ can easily be achieved (e.g. by putting M^+ in a triplet with ν_μ and μ^-) it takes care to arrange that $g_{\nu_\mu M^-} = g_{\nu_\mu \mu^-}$ in models based on the gauge group SU(2)xU(1). For example, if the weak current couples to the doublet

$$\begin{pmatrix} \nu_\mu \\ \mu^-\cos\omega + M^-\sin\omega \end{pmatrix}_L$$

so that $g_{\nu_\mu M^-} = \tan^2\omega\, g_{\nu_\mu \mu^-}$, universality requires $\omega \ll 1$ unless e^- and $(d\cos\theta_c + s\sin\theta_c)$ are similarly mixed with new entities. Interestingly it is possible to arrange that μ^-, e^- and $d(\cos\theta_c + s\sin\theta_c)$ are all mixed by a common amount ω in a natural way[72] (i.e. without specially adjusting the parameters in the Lagrangian; heavy μ's and e's - M^-, M^0, E^-, E^0 - and new quarks - t, b, t', b' - are needed to do this). This has the interesting consequence that the effective coupling of ν_μ to new leptons and u to new quarks is $G_F\tan^2\omega$ which can be bigger than G_F, a possibility worth bearing in mind. Furthermore, since

$$\frac{G_F}{\sqrt{2}} = \frac{g^2\cos^2\omega}{M_W^2}$$

where g is the intrinsic weak coupling constant (which is fixed in terms of e and the Weinberg angle θ_W), the predicted value of M_W is reduced by $\cos\omega$ to

$$M_W = \frac{37.5\cos\omega}{\sin\theta_W}\ \text{GeV}.$$

in SU(2)xU(1). This is a useful reminder that the current theoretical prejudice that $M_W \sim 65$ GeV could be quite wrong. (However, the neutral to charged current ratio is altered by a factor $\sec^2\omega$ in natural models relative to the value it would have in the standard model; since this ratio is successfully predicted, to order 10 or 20%, by the simplest model in which $M_W = M_Z\cos\theta_W$ - see Section 9 - we must either conclude that ω is small or abandon the simplest model in favour of one in which somehow $M_W \approx M_Z\cos\theta_W\cos\omega$).

The problem with universality is partially bypassed in models based on bigger gauge groups (e.g. in an SU(3)xU(1) model with leptons in a triplet) in which the $\nu_\mu \to \mu^-$ and $\nu_\mu \to M^-$ transitions will be mediated by different vector bosons. However, universality now requires that the transition $\nu_\mu \to M^-$ be accompanied by new hadron production at the lower vertex (unless a universal mixing is again introduced) in which case M^- production would have an elevated threshold.

We see that M^- production is very model dependent and that although copious ν_μ production of M^- is possible, it is suppressed in many models. Therefore it is not really possible to conclude that τ (1.9) is not an M^- at present.

9. Neutral Currents

9.1 Introduction

Data on $\nu e \to \nu e$, $\nu p \to \nu p$, $\nu N \to \nu\Delta$, $\nu N \to \nu\pi\ldots$, $\nu N \to \nu\ldots$ and for the corresponding $\bar\nu$ processes now exist. It would neither be appropriate nor possible to give a thorough review of the interpretation of all these data in these lectures. Instead, I shall outline the general strategy and summarize the main conclusions. For the most part, I shall follow the excellent review made last year by Bjorken[6], to which the reader is referred for a more complete discussion.

Two possible approaches are:

a) Confront the data directly with the standard SU(2)xU(1) model and see how long it survives.

b) Try to extract as much information as possible about neutral current couplings from the data in a model independent way.

Both approaches are interesting and instructive and we will use both in turn.

9.2 Neutral Current Couplings in the SU(2)xU(1) Model

In this model there is an SU(2) ("iso") triplet of intermediate vector bosons ($W^\pm_\mu$, W^0_μ) and a singlet B_μ. All left and right handed Fermi fields

$$\psi_L \equiv \frac{(1-\gamma_5)}{2}\psi,\quad \psi_R \equiv \frac{(1+\gamma_5)}{2}\psi$$

must be classified into multiplets of the gauge group. For ν_e and e the only possibility is

$$\psi_L = \begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L,\quad e^-_R$$

given that a $\bar\nu e\, e^-W^+$ coupling is needed. ν_μ and μ^- are classified similarly. The SU(2) invariant interaction Lagrangian is

$$\mathcal{L}_{int.} = \bar\psi_L\gamma^\mu(g\vec{T}\cdot\vec{W}_\mu - g'B_\mu)\psi_L + g''\bar{e}_R\gamma^\mu B_\mu e_R.$$

Whatever the origin of the masses of the vector bosons, W^0_μ and B_μ will mix in general, one eigenstate

$$A_\mu = \sin\theta_W W^0_\mu + \cos\theta_W B^0_\mu$$

being the massless photon and the other

$$Z_\mu = \cos\theta_W W^0_\mu - \sin\theta_W B_\mu$$

being massive.

We can now fix most of the parameters thus

a) The strength of the $\bar{\nu}_e e^- W^+$ vertex gives

$$\frac{g^2}{2M_W^2} = \frac{G_F}{\sqrt{2}} .$$

b) Since ν_e is electrically neutral, the combination $gW^0_\mu - g'B_\mu$ which couples to $\bar{\nu}_e\nu_e$ must be orthogonal to the photon, i.e.

$$\tan\theta_W = g'/g .$$

c) To reproduce the known photon-electron coupling we require

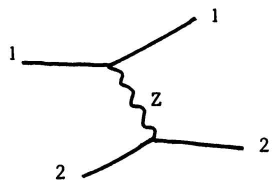
$$e = -2g\sin\theta_W ,\quad g'' = -2g' .$$

At this point there are four parameters in this model (which is due to Glashow[73]): e, G_F, θ_W and M_Z. With a specific mechanism for mass generation, M_Z may be determined theoretically. For example, with the simplest form of the Higgs mechanism (introduced in this context by Weinberg and Salam[74])

$$\eta \equiv \frac{M_W}{M_Z\cos\theta_W} = 1$$

but this is not true in general and we will keep η as a free parameter when making comparisons with the data (if the mass generating mechanism gives $M_{W_3}=M_{W\pm}$ then diagonalization of the mass matrix gives $\eta=1$ whatever the mixing term M_{WB} provided M_B is chosen to make $M_\gamma=0$, as is easy to check). However, we shall see that remarkably the data choose $\eta\simeq 1$.

We can now write down the effective interaction for the process



Generalizing the discussion above to determine the couplings for particles in any SU(2) representation (a useful exercise for the student) the answer is

$$\mathcal{L}^{(1,2)}_{\mathrm{eff.}} = \eta^2\frac{G}{\sqrt{2}}\, 2\left[I_3 - Q\sin^2\theta_W\right]^1_{\lambda_1}\left[I_3 - Q\sin^2\theta_W\right]^2_{\lambda_2}$$
$$\times\ \bar{\psi}_1\gamma_\lambda(1+\lambda_1\gamma_5)\psi_1\,\bar{\psi}_2\gamma_\lambda(1+\lambda_2\gamma_5)\psi_2$$

where $\lambda_i=\pm1$, Q_i is the electric charge and $(I_3)^i_{\lambda_i}$ is the third component of "weak isospin" for particle i with helicity factor $1+\lambda_i\gamma_5$.

It is important to notice that the neutral current coupling is sensitive to the existence of extra leptons and quarks whatever their mass. For example, in the standard mode e^-_R is a singlet and $(I_3)^{e^-}_R=0$ but it could belong to a doublet with an as yet undiscovered heavy neutral lepton

$$\begin{pmatrix} E^0 \\ e^- \end{pmatrix}_R$$

in which case $(I_3)^{e^-}_R$ would be $-1/2$ and the $\bar{e}eZ$ coupling would be different (in particular it would have no axial part).

We will see how the SU(2)xU(1) model compares with the data at the end of the next section.

9.3 General Phenomenology

We write the effective interaction which determines $\sigma(\nu N\to\nu x)$ in the form

$$\mathcal{L}_{\mathrm{eff.}} = \frac{G}{\sqrt{2}}\,\bar{\nu}_\mu\gamma_\lambda(1-\gamma_5)\nu_\mu$$
$$\times \sum_i\left\{\epsilon_L(i)\bar{\psi}_i\gamma_\lambda(1-\gamma_5)\psi_i + \epsilon_R(i)\bar{\psi}_i\gamma_\lambda(1+\gamma_5)\psi_i\right\}$$

where i runs over all fermions. Note that:

a) We have assumed that there are V and A but not S, P, T interactions (pure S and/or P are ruled out by the fact that the y distribution in $\nu N\to\nu$... is not y^2).

b) We have assumed that the interaction is diagonal in lepton flavours (e.g. $\nu\to\nu$ but there is no new lepton contribution $\nu\to\nu'$) and quark flavours (e.g. $d\not\to s$; $u\not\to c$). $u\to c$ transitions would have interesting consequences for neutrino reactions, which the reader can work out, but they are severely limited by the fact that large D^0-$\bar{D}^0$ mixing has been excluded by e^+e^- experiments[45]. We will therefore not discuss flavour changing neutral currents further except to mention that the plausible requirement that they be eliminated "naturally" (i.e. without specially adjusting the parameters in the Lagrangian) puts strong constraints on possible models[75].

c) We have not included a $\bar{\nu}\gamma_\lambda(1+\gamma_5)\nu$ piece because it has no effect in experiments using left handed ν_μ's derived from π and K decay. However, its possible presence should not be forgotten in drawing general conclusions e.g. contrary to some statements in the literature, it is possible to have a parity conserving Lagrangian but still have $\epsilon_L(i)\neq\epsilon_R(i)$, so that the observations that $\sigma(\nu e\to\nu e)\neq\sigma(\bar{\nu}e\to\bar{\nu}e)$ and $\sigma(\nu N\to\nu+\ldots)\neq\sigma(\bar{\nu}N\to\bar{\nu}+\ldots)$ are not in themselves evidence for parity violation (neither is $\sigma(pp)\neq\sigma(\bar{p}p)$ evidence for parity violation in the strong interactions!)

We now summarize briefly some of the conclusions which can be drawn from the data and some of the evidence for them:

i) $\epsilon_L(\text{quark})\neq\epsilon_R(\text{quark})\neq0$. Some evidence[76]:

a) $\dfrac{\sigma_{\bar{\nu}p\to\bar{\nu}p}}{\sigma_{\nu p\to\nu p}} = 0.4\pm0.2 \neq 1 \to \epsilon_L\neq\epsilon_R$.

b) Deep inelastic data. Defining

$$|\epsilon_{L,R}|^2 = |\epsilon_{L,R}(u)|^2 + |\epsilon_{L,R}(d)|^2$$

the naive parton model (cf. Section 4.3) gives

$$\frac{d\sigma^{\nu N\to\nu X}}{dx\,dy} = \frac{G^2 s}{2\pi}\Big[xQ|\epsilon_L|^2 + x\bar{Q}|\epsilon_R|^2 + (1-y)^2\{x\bar{Q}|\epsilon_L|^2 + xQ|\epsilon_R|^2\}\Big]$$

$$\frac{d\sigma^{\bar{\nu}N\to\bar{\nu}X}}{dx\,dy} = \frac{G^2 s}{2\pi}\Big[\,-\,-\,\epsilon_L\leftrightarrow\epsilon_R\,-\,-\,-\,\Big]$$

for an I=0 nucleus if we neglect all but u, d, $\bar{u}$ and $\bar{d}$ quarks. Fits to the y distributions measured in the CITF experiment[76] parametrised as

$$\frac{d\sigma^{\nu,\bar{\nu}}_{N.C.}}{dy} = \frac{G^2 s}{2\pi}\left[g_{L,R} + (1-y)^2 g_{R,L}\right]$$

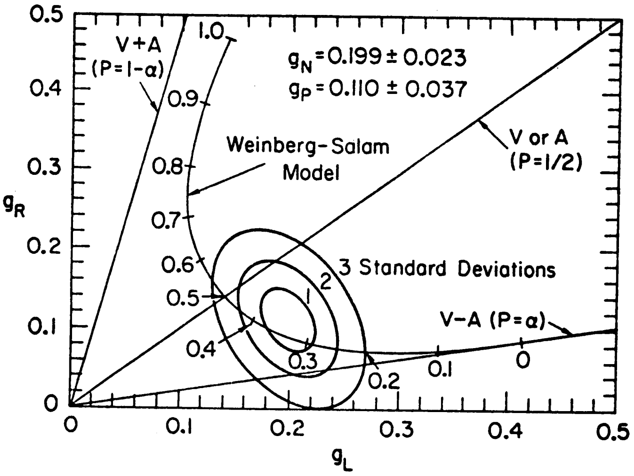


Fig. 15. The coefficients g_R and g_L defined in the text. The numbers on the Weinberg-Salam curves correspond to different values of $\sin^2\theta_W$.

are shown in Fig. 15 ($g_N \equiv g_L$, $g_P \equiv g_R$, in the experimenters' notation). The V±A lines assume the rather large value

$$\alpha \equiv \frac{\int x\bar{Q}\,dx}{\int x(Q+\bar{Q})\,dx} = 0.17$$

(with a smaller value they would lie closer to the axes). We conclude that $\varepsilon_L \neq 0$ and (at the two standard deviation levels) $|\varepsilon_L| \neq |\varepsilon_R|$, $\varepsilon_R \neq 0$.

Using $\frac{\sigma^{\bar\nu}}{\sigma^{\nu}} = 0.38\pm0.2$ to fix $\alpha = 0.05\pm0.02$, the total cross-section measurements give[6)]

$$|\varepsilon_L|^2 = 0.26\pm0.04$$

$$|\varepsilon_R|^2 = 0.06\pm0.05$$

(where the error assignments should not be taken seriously, especially for ε_R). A note of caution is necessary. The CITF experiment is at an energy at which we have seen that scaling violations are important according to asymptotic freedom and a naive parton model approach may be misleading. It will be important to bear this in mind in interpreting the much more accurate data which will soon be available from SPS experiments.

ii) $\varepsilon(u) \neq \pm\varepsilon(d)$, i.e. there are both I=0 and I=1 contributions. Some evidence:

a) On an I=0 target an I=0 neutral current would give

$$\sigma(\nu\to\nu\pi^+..) = \sigma(\nu\to\nu\pi^-..) = \sigma(\nu\to\nu\pi^0..)$$

and likewise for $\bar\nu$. The Gargamelle data give

$$\pi^0 : \pi^+ = 1.4\pm0.2 \qquad \text{for } \nu$$

$$= 2.1\pm0.4 \qquad \text{for } \bar\nu$$

The number of neutrons and protons in Gargamelle are not in fact quite equal (i.e. $I\neq0$) but the prediction is not much changed by this, i.e. a pure I=0 current is excluded.

b) $$\frac{\nu n\to\nu n\pi^+\pi^-}{\nu p\to\nu p\pi^-\pi^+} = 0.49\pm0.20 \neq 1$$

This requires interference between I=0 and I=1 currents i.e. both are needed.

We see that a lot of model independent information about the ε's can be derived from existing neutrino data (it is gratifying that different measurements yield consistent results[6)]). The results are all consistent with the standard SU(2)xU(1) model (with ν_e, e, ν_μ, μ, u, d, s and c) with $\eta\simeq1$ and $0.2<\sin^2\theta_W<0.4$, as the reader can easily check in the case of the results quoted here. However, there are experiments which look for parity violating effects in atoms which do not seem to agree with the simple model[77)]. A large atom (Bismuth) is used so the vector coupling to the nucleus (which is coherent between different nucleons) dominates and the experiments measure $g_A^e \times g_V^{Bi}$. No parity violation has been seen at a level about ten times smaller than expected according to the standard model, combined with atomic calculations. There is still some small doubt about the validity of the atomic calculations but assuming they are correct we must conclude that:

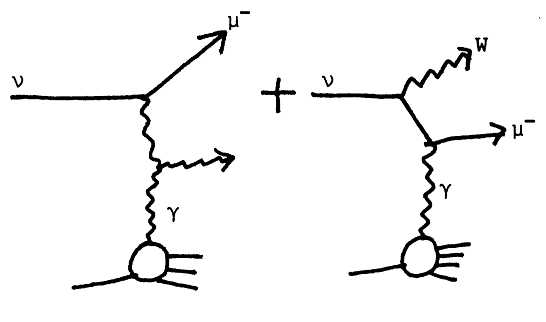
either $g_A^e\simeq0$ which requires $(I_3)_R^{e^-}\simeq(I_3)_L^{e^-}$ (approximate equality could occur in models with mixing); the simplest possibility is that e_R^- forms a doublet with a heavy E^0, as discussed above;

or $g_V^{Bi}\simeq0$. By a judicious but artificial mixture of I=0 and I=1 currents we could make $g_V^{Bi}=0$. However, it is probably sufficient to make a model with a pure I=1 current (so that $g_V^{(A,Z)} \sim \frac{A-2Z}{A}$). In contrast to models with $g_A^e=0$, we would expect to find parity violation in hydrogen in this case.

Clearly atomic physics experiments can give extremely important information and the outcome of experiments using Caesium[78)] (for which the atomic calculations are presumably very reliable) is eagerly awaited. Further information on the coupling of neutral currents to charged leptons will come from $e^+e^-\to\mu^+\mu^-$, $e^+e^-\to e^+e^-$ and $e/\mu p\to e/\mu+\ldots$ experiments.

10. Finding the W

A mechanism for W production by neutrinos is



A signature is provided by the decay $W^+\to\mu^+\nu$ which gives a large p_T μ^+. Failure to observe this process has been used to set the limit $M_W>8$ GeV[79)]. Unfortunately it requires $E_\nu>5000$ GeV to have a chance of finding the W in this way if $M_W\simeq65$ GeV.[80)]

Virtual W's give a characteristic $\left(\frac{1}{1+Q^2/M_W^2}\right)^2$ scaling violation in neutrino reactions which might be used to set limits of 20 GeV or so on M_W in current neutrino experiments. This propagator effect would show up very clearly in measurements of $\sigma(e^-p\to\nu\ldots)$ which could be made with the s=27,000 GeV^2

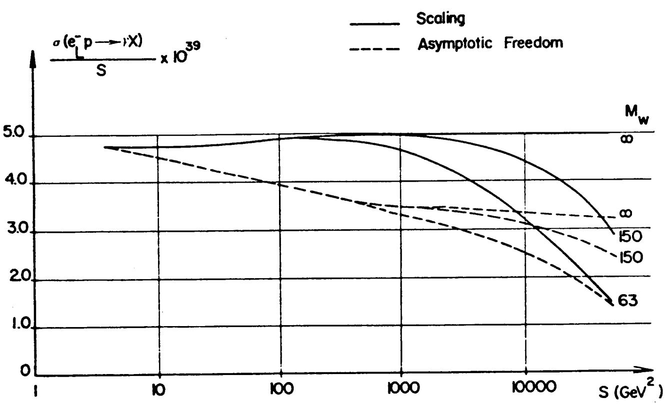


Fig. 16. The cross-section for left-handed electrons divided by s as a function of s in various cases[81].

ep storage ring CHEEP under discussion at CERN[82], as can be seen from Fig. 16. In fact measurements of $e^- p \to \nu$... with this device may provide the next step forward in energy for neutrino physics[81]; they are equivalent to experiments with a 15,000 GeV neutrino beam! Unfortunately, however, ep rings are not good devices for directly producing W's or Z^o's. It seems most likely that this will first be done with very large pp or $p\bar{p}$ storage rings[83] (although e^+e^- storage rings provide the best environment for studying W's and Z^o's[84]).

11. Concluding Remarks

Neutrino experiments are clearly an extremely powerful way to investigate the weak interactions and probe the structure of hadrons. For example, they can measure the quark distributions and test whether the strong interactions are really an asymptotically free gauge theory. They can search for new leptons, new quarks and new currents and determine the neutral current interactions.

Recent progress in neutrino physics is illustrated in the accompanying chart. In the left hand column I have listed some questions about neutrinos raised more than seventeen years ago by Lee and Yang[85], when they spelled out the theoretical reasons for doing neutrino experiments, together with some others added by me[86] when I reviewed the situation in 1972. Comparison of the answers known in 1972 and mid-1977 shows that we have made spectacular progress in the last five years. I hope and believe that progress in the next few years will be equally exciting.

Some History

Questions	Answers 1972	Answers 1977
$\nu_\mu = \nu_e$?	$\nu_\mu \neq \nu_e$	
Lepton Cons. L^+?	$\frac{\nu \to \mu^+}{\nu \to \mu^-} < 5\%$	$\frac{\nu \to \mu^+}{\nu \to \mu^-} < \frac{1}{2}\%$
Neutral Currents?	Poor limits	Yes
'Locality' Vector nature of W.I.		Consistent
Universality between ν_μ and ν_e?		$\frac{\nu_e \to e..}{\nu_\mu \to \mu..} = 1.26 \pm .23$
Charge Symmetry?		Apparent viol. due to Charm?
CVC; Isotriplet Current		In $\nu N \to \mu \Delta$ $-.51 \pm .11 < A_2/A_1 < .23 \pm .07$
W?	$M_W > 1.8$ GeV	$M_W > 8$ GeV
What happens as $E_\nu \to$ 'unitarity limit'?		
Heavy leptons?		Yes-but wit own lepton no.? (only seen in $\bar{e}e$)
Charm?		Yes
Scaling, sum rules, etc.?		Much information

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