Ref.TH.3057-CERN

A CONFINING MODEL OF THE WEAK INTERACTIONS

L.F. Abbott^{*)}

and

E. Farhi

CERN--Geneva

ABSTRACT

The weak interactions are analyzed as the low energy effective interactions of a strong coupling gauge theory whose scale is $G_F^{-1/2}$. The light fermions are shown to be bound states obeying 't Hooft's consistency conditions. The symmetries of the theory are used to analyze the low-energy interactions. Instanton mediated baryon violations are discussed and experimental signatures at high energy are presented.

---

^{*)} Permanent address : Brandeis University, Waltham, MA 02254, U.S.A.

Ref.TH.3057-CERN

7 March 1981

1. INTRODUCTION

At low energies the weak interactions are well described by the usual charged and neutral current four Fermi Lagrangian with coupling strength G_F ($G_F^{-1/2} \approx 300$ GeV). In the standard $SU(2)_L \times U(1)$ model¹⁾ these interactions arise from the exchange of weakly coupled massive vector bosons and the mass scale for G_F is set by the vacuum expectation value of a scalar field. However, only the low-energy behaviour of the weak interactions has at present been determined, so there exists another possibility. Perhaps the four Fermi interactions we see are instead the low-energy remnants of a strong interaction between fermions which are bound states of confined constituents. In this case, G_F is determined by the mass scale of this underlying strong interaction, i.e., the Λ parameter of the theory. Thus, we imagine that the weak interactions are a result of a confining gauge theory²⁾ which binds constituents into the observed fermions at a scale of about $G_F^{-1/2}$.

In Ref. 2), a model based on this idea was presented. This model is very similar to the standard $SU(2)_L \times U(1)$ theory, except that the $SU(2)_L$ gauge theory is taken to be confining rather than spontaneously broken. It was shown that this model has light bound states corresponding to the observed fermions and produces the correct four Fermi weak interactions. This model may be a viable alternative to the standard spontaneously broken theory. Like the standard theory, the confining model is based on a renormalizable Lagrangian and is thus well defined.

In this paper we will address certain questions which were not answered in our previous work. We begin by reviewing the basic assumptions of the model and by discussing its symmetry properties. Using these symmetry properties, we will show that the massless bound states appearing in the theory satisfy the 't Hooft consistency conditions³⁾ for the existence of massless bound states in confining theories. Again, using the symmetries we will show that the low energy effective Lagrangian, which correctly describes the weak interactions, is included in the four Fermi interactions of this theory and that only one other possible interaction is allowed. We will discuss why this extra term might be small. Then we address the question of whether a strong coupling $SU(2)_L$ theory might have unacceptably high levels of baryon number violation through instantons. The answer is no and this is shown in Section 5. Finally, we examine some of the experimental signatures of strong coupling weak interactions.

Although our approach has much success, there are certain difficulties. The basic mass scale associated with the weak interactions is $G_F^{-1/2} \approx 300$ GeV. In order to account for neutral current interactions, the left-handed fermions must have electromagnetic form factors which at low Q^2 go like Q^2/M_1^2 where

$M_1 \approx 100$ GeV. On the other hand, a bound state structure for the muon will produce an anomalous magnetic moment of order m_μ^2/M_2^2 and from experiment we know that $M_2 \geq 600$ GeV. The discrepancy between these mass scales is unexplained.

Our entire analysis of the weak interactions depends very little on the nature of the confining force. There are three basic features which a confining gauge theory must have in order to provide a possible description of the low-energy weak interactions. First, it must have light (on the scale $G_F^{-1/2}$) composite fermions. Second, it must have a global $SU(2)$ symmetry under which the bound state fermions are doublets. Third, the strong interaction must only act on left-handed physical fields. These last two conditions almost guarantee the correct form for the low-energy effective Lagrangian.

All of these conditions are met by the $SU(2)_L$ confining model, but perhaps it is not unique in this regard. In particular, since there is no spontaneous symmetry breaking in our approach, there is no fundamental reason for a scalar field to be present. In the $SU(2)_L \times U(1)$ model, a scalar doublet is necessary to bind with fermionic preons to produce the observed fermions. However, if there exists a confining theory based on a different group which has the necessary three features, perhaps a weak interaction model with no elementary scalars can be constructed. We are presently searching for such a model.

2. THE MODEL AND ITS SYMMETRIES

In this section we will review some of the basic features of the strong-coupling weak interaction model introduced in Ref. 2). This model is based on the same Lagrangian as the standard $SU(2)_L \times U(1)$ model. For three generations of fermions, there are 12 (one lepton and three coloured quarks for each generation) left-handed $SU(2)_L$ doublet Fermi fields ψ_{Li}^a ($a=1-12$; $i=1-2$) and a complex $SU(2)_L$ doublet scalar field ϕ_i . In addition, there are 24 right-handed Fermi fields U_R^a, D_R^a which do not feel the $SU(2)_L$ force. These particles have the same $U(1)$ quantum number assignments as in the standard model. They interact through $SU(2)_L$ and $U(1)$ gauge interactions, Yukawa couplings and scalar self couplings. The only changes which are made in the strong coupling model are in the values of certain parameters in the Lagrangian. These changes are:

- 1) The potential of the scalar field is modified so that the scalar vacuum expectation value $\langle \phi \rangle$ vanishes. Therefore, there is no spontaneous symmetry breaking.
- 2) The $U(1)$ coupling constant is set equal to e , the electromagnetic coupling constant where $e^2/4\pi = 1/137$.

- 3) The $SU(2)_L$ coupling constant is increased in value so that the $SU(2)_L$ interactions become strong at an energy scale of order $G_F^{-1/2}$.

Since the $SU(2)_L$ gauge symmetry is unbroken, and this theory becomes strong coupling, confinement presumably occurs and all physical states must be $SU(2)_L$ singlet bound states. Most of these bound states will have masses of order $G_F^{-1/2}$ since this is the mass scale determined by the running $SU(2)_L$ coupling constant. However, this model contains physical fermions which have masses much lighter than $G_F^{-1/2}$. In Ref. 2) this was demonstrated by explicitly identifying the chiral symmetry responsible for keeping the fermions light. It was also shown that the $U(1)$ quantum numbers of these light fermions agree with the usual charge assignments^{3),4)}. Since the $U(1)$ is unbroken and it has coupling constant e , the physical states have the correct electromagnetic interactions.

In the next section we will rederive the light fermion spectrum using the consistency conditions³⁾ of 't Hooft. In Section 4 we will discuss the effective weak interactions of the light fermions. For these discussions it is useful to identify the global symmetries of the confining $SU(2)_L$ Lagrangian.

The $SU(2)_L$ confinement occurs at an energy scale of about $G_F^{-1/2}$. The $SU(3)$ colour interactions, electromagnetic interactions and Yukawa couplings are all weak compared to the $SU(2)_L$ interactions at the scale $G_F^{-1/2}$ so these interactions can be neglected. In this case there is no distinction between quarks and leptons or between particles in different generations, so the 12 left-handed $SU(2)_L$ doublets ψ_{Li}^a have an $SU(12)$ symmetry. (We are not interested in the symmetries of the right-handed fields since they do not feel the strong force.) Naively this symmetry would be $U(12)$ but the $U(1)$ subgroup has an anomaly in the presence of $SU(2)_L$ interactions. The effects of this anomaly will be discussed in Section 5.

In addition to the $SU(12)$ symmetry of the fermions, there is a global $SU(2)$ symmetry associated with the scalar doublet. If we write the scalar doublet in a matrix form

$$\Omega = \begin{pmatrix} \phi_1 & -\phi_2^* \\ \phi_2 & \phi_1^* \end{pmatrix} \quad (2.1)$$

then the $SU(2)_L$ standard Lagrangian is invariant under the transformation

$$\Omega \rightarrow \Omega u \quad (2.2)$$

where u is an arbitrary $SU(2)$ matrix. [This is the same symmetry discussed in Ref. 2) up to a global gauge transformation.]. This symmetry is only exact in the absence of electromagnetism and Yukawa couplings. Up to these small corrections, the global symmetry of the strong $SU(2)_L$ theory is $SU(12) \times SU(2)$.

3. LIGHT COMPOSITE FERMIONS

't Hooft has given the consistency conditions³⁾ for the existence of light fermion bound states in a confining theory. We will now show that the light fermions in our $SU(2)_L$ strong theory satisfy these conditions.

't Hooft's analysis begins by identifying the global symmetries of a given model. The elementary, confined fermions then produce certain anomalies in these global symmetries through the usual triangle diagrams. We next postulate the existence of certain massless composite fermions, and calculate the anomalies in the global symmetry group produced by these massless states. In order for the postulated existence of these massless composites to be consistent, 't Hooft requires that the anomalies produced by the fundamental fermions be identical to those of the bound state massless fermions.

As we have seen, the global symmetry of the $SU(2)_L$ model is $SU(12) \times SU(2)$. The only possible anomaly can come from triangle diagrams with three $SU(12)$ generators at the vertices. If we normalize the anomaly so that a left-handed fermion in the fundamental representation of $SU(12)$ produces an anomaly of one unit, then the total anomaly of the underlying theory is two since the ψ_{Li}^a are two 12's of $SU(12)$. This value must be matched by the anomaly of the bound state massless composites.

The composite fermions consist of bound states of the fundamental fermion doublets and the scalar field doublet. For each fermion type, we can construct two different $SU(2)_L$ singlets with the scalar field. In fact, these two gauge singlets will be a doublet of the global $SU(2)$. Explicitly:

$$\Psi_L^a = \begin{pmatrix} \phi_i^* \psi_{Li}^a \\ \epsilon_{ij} \phi_i \psi_{Lj}^a \end{pmatrix} = \Omega^T \psi_L^a \quad (3.1)$$

where the ψ_L^a are the fields of the physical fermions. Note that ψ_L^a is a doublet under $SU(2)_{\text{global}}$ but a singlet under $SU(2)_L$ while ψ_{Li}^a is a global singlet and a doublet under the gauge group. Therefore, the bound state fermions ψ_L^a are two 12's of $SU(12)$ and the anomaly of the composite equals the anomaly in the underlying theory. Thus we have reason to believe that in the limit where $SU(12) \times SU(2)_{\text{global}}$ is an exact symmetry, the composites we have identified should be truly massless.

The colour and electromagnetic interactions violate the $SU(12) \times SU(2)_{\text{global}}$ symmetry. The $SU(3) \times U(1)$ group of these interactions is a subgroup of $SU(12) \times SU(2)_{\text{global}}$ as far as left-handed fields are concerned. Even in the presence of these interactions, there is enough chiral symmetry [before $SU(3)$ dynamical chiral symmetry breaking] to guarantee the masslessness of the states we considered. However, when Yukawa couplings are included, this chiral symmetry is explicitly broken and the composite fermions become massive. Since the masses must vanish when the Yukawa couplings are turned off, the masses will be of order $\lambda G_F^{-1/2}$ where λ is the appropriate dimensionless Yukawa coupling. Typically, λ is much less than one and the physical fermions have masses much less than the scale $G_F^{-1/2}$.

4. FOUR FERMI INTERACTIONS

In our previous paper²⁾ we showed that the strong coupling $SU(2)_L$ theory can give rise to an effective four Fermi interaction which describes low-energy weak interaction physics. We now want to show that this interaction is the only possible type of interaction consistent with the symmetries of the theory except for one other possible term. We will then discuss why this extra term could be small. In fact, the observation of this extra interaction at low energy could signal the existence of a confined theory of the weak interactions⁵⁾.

At energies well below $G_F^{-1/2}$, four Fermi interactions between the bound state fermions are induced by the strong $SU(2)_L$ force. The scale of the $SU(2)_L$ theory is set by requiring that the coefficient of these four Fermi interactions is G_F which is why we say that the $SU(2)_L$ theory becomes strong at an energy of order $G_F^{-1/2}$. Up to small corrections due to colour, electromagnetic and Yukawa couplings, these low-energy four Fermi terms must respect the global $SU(12) \times SU(2)$ symmetry of the underlying $SU(2)_L$ theory. Furthermore, the interactions can only involve left-handed Fermi fields since only these feel the $SU(2)_L$ force. This condition, along with the $SU(12) \times SU(2)$ symmetry greatly restricts the possible four Fermi terms which can appear.

The interactions involve the physical fermions Ψ_L^a which are a (12, 2) under $SU(12) \times SU(2)_{\text{global}}$. The four Fermi interaction must be a singlet under $SU(12)$ so it must contain two Ψ_L^a 's and two $\bar{\Psi}_L^a$'s. Writing the interaction in the order $\bar{\Psi}\Psi\bar{\Psi}\Psi$, we see that the left-handed nature of the Fermi fields forces the interaction to be (V-A) $\times$ (V-A). $\bar{\Psi}\Psi$ decomposes into (1, 1) + (1, 3) + (143, 1) + (143, 3) under $SU(12) \times SU(2)$. However, the last two representations can be Fierz transformed into the first two so the most general low-energy Lagrangian is⁵⁾:

$$\frac{4G_F}{\sqrt{2}} \left[\vec{J}_{\mu L} \cdot \vec{J}_L^\mu + \xi J_{\mu L}^0 J_L^{0\mu} \right] \quad (4.1)$$

where

$$\vec{J}_{\mu L} = \sum_{a=1}^{12} \bar{\Psi}_L^a \gamma_\mu \vec{T} \Psi_L^a \quad (4.2)$$

and

$$J_{\mu L}^0 = \sum_{a=1}^{12} \bar{\Psi}_L^a \gamma_\mu \mathbb{1} \Psi_L^a \quad (4.3)$$

The first term in Eq. (4.1) gives the correct charged-current interaction and part of the neutral current, so the coefficient must be $4G_F/\sqrt{2}$. Note that the general $SU(2)$ structure of these interactions was generated by the global $SU(2)$ symmetry. The universality of the weak interactions is guaranteed by the global $SU(12)$ symmetry.

On the other hand, the interactions of the second term of Eq. (4.1) are not observed and it is known⁵⁾ from neutral current experiments that $|\xi| \leq 0.05$. Thus, we must discuss why ξ is so small in our model and also identify the complete neutral current.

The first term of Eq. (4.1) can be produced by the exchange of a spin one boson of mass $\approx G_F^{-1/2}$ transforming like a triplet under the $SU(2)$ global symmetry group. This state can be made by binding together two scalars into an $SU(2)_L$ singlet state. The lowest lying scalar-scalar bound states⁴⁾ are thus identified with the W^+ , W^- and W^0 and give the leading contribution to the first term of Eq. (4.1). If we try to construct the scalar-scalar spin one bound state which is the $SU(2)$ global singlet partner of the $W^+W^-W^0$, we find that

it is forbidden by Bose symmetry. Thus, as far as scalar-scalar bound state exchange is concerned, the first term in Eq. (4.1) can arise, but the second one cannot.

However, we can also construct a spin one bound state out of two fundamental fermions. Since the fermions are global $SU(2)$ singlets, the bound state is also a singlet and can contribute to the second term of Eq. (4.1), but not to the first. Therefore, we must assume that this state is somewhat decoupled. For example, if it is twice as massive as the W^0 and has one half the corresponding coupling to fermions, we will have $|\xi| < 0.07$.

It will take a more detailed understanding of dynamics to know which states really contribute. However, there is a possibility worth entertaining. If the $SU(2)_L$ interactions lead to a confining force which is spin independent, then it may be impossible to make a massive bound state out of massless fermions without chiral symmetry breakdown⁶⁾. We have been assuming all along that $SU(2)_L$ does not break chiral symmetry, otherwise our entire analysis is spoiled. Therefore, it may be impossible to form the massive fermion-fermion bound state responsible for the second term of Eq. (4.1). Of course this is only speculation.

As was shown in Ref. 2), the complete neutral current interaction can be understood in this strong coupling model. Because the left-handed fermions are bound states, they will have non-trivial electromagnetic form factors at an energy scale of $G_F^{-1/2}$. These form factors produce the remaining terms in the neutral current interaction^{5),7)} not present in Eq. (4.1). This can be seen by assuming that, as in QCD, these form factors can be computed through vector-meson dominance⁷⁾. We assume that the photon can turn into a W^0 boson with amplitude XQ^2 . Then, if X is proportional to the usual $\sin^2\theta$, the neutral current interaction of this model agrees exactly with that of the standard model and thus with the experimental data^{2),5),7)}. If we put in the expected values of the couplings [see Ref. 2)], we discover that X is about ten times bigger than the corresponding number describing $\gamma - \rho^0$ mixing in QCD. This may be understandable since the W^0 is two bosons not two fermions and the binding force is $SU(2)$ not $SU(3)$. Still the largeness of X may indicate that the assumption of vector meson dominance is not warranted in this model. With or without vector meson dominance, the form factors required to produce the observed neutral current have a scale only slightly less than $G_F^{-1/2}$ and since they arise from binding at $G_F^{-1/2}$ this seems perfectly consistent.

5. BARYON NUMBER VIOLATION THROUGH INSTANTONS

Baryon number is not conserved in the $SU(2)_L \times U(1)$ model because of $SU(2)_L$ instantons⁸⁾. This is true both for strong and weak couplings. The calculation of the amplitude for a baryon number violating process yields a factor

of $\exp(-8\pi^2/g_2^2)$ which, for the standard weak coupling theory ($g_2 \approx 0.6$), is a tremendous suppression factor. Baryon number violation occurs, but at an unobservable rate. However, for the strong coupling theory, g_2 gets large and the suppression factor is not operative⁹⁾. Baryon number violation may occur at an appreciable rate. We will now estimate this rate and find that it is still orders of magnitude smaller than the experimental bound.

At the scale $G_F^{-1/2}$ where $SU(2)_L$ becomes strong, we can ignore colour and electromagnetism. We will assume that there are three generations of quarks and leptons, so at the scale $G_F^{-1/2}$ there are twelve indistinguishable $SU(2)_L$ doublets ψ^a ($a=1, 12$) as in Section 2. The conservation of fermion number for a fermion doublet of type a , N_a , is violated by $SU(2)_L$ instantons, but the difference of any two, $N_a - N_b$, is always conserved. The lowest dimension operator which can mediate baryon number violation, conserves the various $N_a - N_b$ and is $SU(12)$ symmetric is⁸⁾:

$$O = \epsilon_{a_1 a_2 \dots a_{12}} \psi^{a_1} \psi^{a_2} \dots \psi^{a_{12}} \quad (5.1)$$

where the $SU(2)$ indices are tied up in all possible ways. The amplitudes for baryon number violating processes are given by the appropriate matrix elements of an effective Hamiltonian involving O . The only relevant dimensional parameter here is $G_F^{-1/2}$ so on dimensional grounds this effective Hamiltonian must go like

$$H \approx G_F^7 O \quad (5.2)$$

The effective Hamiltonian of Eq. (5.2) mediates processes with $\Delta B = 3$. One such process is nucleon + nucleon $\rightarrow$ antinucleon + 3 antileptons. For this process, all the mass scales associated with the matrix element of (5.2) and the integrations over final state momenta are of order m_{proton} or less. Of the 12 fields appearing in O , only three quark fields belong to the first generation. Three more quark fields belong to the second generation so for each of them there is a suppression of $\sin\theta_c$ for a matrix element involving ordinary nucleons. The last three quark fields belong to the third generation and so are suppressed by Kobayashi-Maskawa angle factors. We will assume that each of them is doubly Cabibbo suppressed by a factor $\sin^2\theta_c$. Putting all this together, we find that the rate for nucleon + nucleon $\rightarrow$ antinucleon should be about

$$R \approx (\sin \theta_c)^{18} G_F^{14} m_{\text{proton}}^{29} \approx (2 \times 10^{50} \text{ years})^{-1} \quad (5.3)$$

The factor of $(m_{\text{proton}})^{29}$ is an upper limit since nuclear physics effects bring factors even smaller than m_{proton} . Still this rate is many orders of magnitude smaller than present limits on the photon lifetime and is unobservable.

A process in which two nucleons go into a state of three strange quarks like the Ω will be somewhat larger than (5.3) because it involves fewer Cabibbo suppression factors. The rate for this process is roughly

$$R \approx (\sin \theta_c)^{12} G_F^{14} m_{\text{proton}}^{29} \approx (2 \times 10^{46} \text{ years})^{-1} \quad (5.4)$$

which is still unobservable. However, there is considerable uncertainty in the estimation of these rates. For example, if the factor G_F in Eq. (5.4) is replaced by $(1/100 \text{ GeV})^2$ we find a lifetime of about 2×10^{33} years. This is not too far from present bounds on the proton lifetime. Thus, someday, baryon violation from $SU(2)_L$ interactions, if they are strong, may be observable.

6. HIGH-ENERGY PHENOMENOLOGY

The high-energy properties of the $SU(2)_L$ confining theory of the weak interactions are very different from those of the standard model. At energies of the order of $G_F^{-1/2}$ the strong nature of the weak interactions should clearly reveal itself. In this section we discuss some of the distinctive features of this high-energy phenomenology.

Although the conventional $W^\pm$ and Z^0 will not be seen if the confining version of the weak interactions is correct, bound-state vector bosons with similar properties will appear. If we assume that the low-energy weak interactions are dominated by the exchange of the lowest lying bound-state vector bosons, then we find an upper limit for the mass of the bound state $W^\pm$ which was first derived by Bjorken⁵⁾ and by Hung and Sakurai⁷⁾. This bound requires that the $W^\pm$ in the confining theory have a mass less than 166 GeV. On the other hand, we would expect²⁾ the bound-state $W^\pm$ to be more massive than the $W^\pm$ of the

standard model. Thus, the $W^\pm$ mass should be somewhere between 100 GeV and 166 GeV. The W^0 mass should be somewhat heavier⁷⁾ than the $W^\pm$ mass as in the standard model. The $W^\pm$ and W^0 are quite analogous to the $\rho^\pm$ and ρ^0 of ordinary hadronic physics. The width of the ρ is about 20% of its mass. By analogy then, we expect the $W^\pm$ and W^0 to be broad resonances, around 10 - 30 GeV in width.

The ordinary $W^\pm$ and Z preferentially decay into single fermion pairs. The $W^\pm$ and W^0 of the strong coupling theory will also have these channels open, but in addition they will strongly decay into many fermion final states. At the scale of the mass of the heavy bosons, quarks and leptons of all generations are essentially equivalent so we expect equal numbers of leptons and quarks of each colour divided uniformly between the generations in these final states. The large numbers of leptons produced per event is a particularly good signal¹⁰⁾ of the strong nature of the weak interactions.

The new strong interaction sector will, of course, have a rich spectrum like the hadronic spectrum of QCD. However, since the mass scale of this strong interaction is of order $G_F^{-1/2}$, most of the excited states will lie about 300 GeV above the $W^\pm$ and W^0 and so will not be easily detectable. However, in addition to the $W^\pm$ and W^0 , there should exist the first excited states of the ordinary fermions, the Higgs-like scalar bound-state and perhaps a ω like bound-state in the region 100 - 200 GeV. Like the W 's, these states will be broad and can decay into numerous particles and in particular produce large numbers of leptons.

We have based many of our arguments about the strongly coupled theory responsible for the weak interactions on analogous properties of QCD and hadronic physics. However, the strong interactions we are describing differ in two important respects. One is that they are purely left-handed and thus we expect parity violation in strong interaction processes¹¹⁾. The other is that there is no chiral symmetry breaking in these strong interactions. The effects of this on our analogies with QCD are not known.

7. CONCLUSIONS

In this paper we have continued our analysis of a model of the weak interactions without spontaneous breakdown. We have strengthened our arguments for the existence of light bound state fermions. Using the global symmetries of the theory we have shown how to derive the correct form for the low-energy weak interaction Lagrangian including universality. We also discussed other possible interactions which can arise at low energy, at high energy and through instanton effects

There are two difficulties remaining. The first is to understand why the electric form factors required to give the neutral current are much larger than the magnetic form factors. The second is to better understand the suppression of the isosinglet channel in the neutral current.

Our theory has many phenomenological successes. It suggests the possibility of constructing other models without fundamental scalars based on confining gauge theories. Proton-anti-proton experiments should soon determine whether or not the weak interactions are strong.

ACKNOWLEDGMENTS

We are grateful to many of our colleagues at CERN for their helpful comments, in particular A. De Rújula, J. Ellis and R. Barbieri. Some of the ideas presented here were analyzed independently by T. Banks, Y. Frishman, E. Rabinovici and S. Yankielowicz and we have benefited from discussions with them. This work was completed at the Weizmann Institute. We thank H. Harari and the members of the Theory Group for their hospitality.

REFERENCES

- 1) S. Weinberg, Phys. Rev. Lett. 19 (1967) 1264;
A. Salam in Elementary Particle Theory : Relativistic Groups and Analyticity
(Nobel Symposium No. 8) ed. N. Svartholm (Almqvist and Wiksell, Stockholm,
1968).
- 2) L. Abbott and E. Farhi, CERN preprint-TH 3015 (1981) - Phys. Lett. to be
published.
- 3) G. 't Hooft, Cargèse Summer Institute Lectures (1979), Plenum Press, New York
(1980).
- 4) T. Banks and E. Rabinovici, Nucl. Phys. B160 (1979) 349;
E. Fradkin and S.H. Shenker, Phys. Rev. D19 (1979) 3682;
S. Dimopoulos, S. Raby and L. Susskind, Nucl. Phys. B173 (1980) 208.
- 5) J.D. Bjorken, Arctic Summer School, Akaslompolo, Finland (1980); Phys. Rev.
D19 (1979) 335.
- 6) A. Casher, Phys. Lett. 83B (1979) 141.
- 7) P.Q. Hung and J.J. Sakurai, Nucl. Phys. B143 (1978) 81.
See also N. Dombey, Neutrino 79, Vol. II, ed. A. Haatuft and C. Jarlskog,
Bergen (1979).
- 8) G. 't Hooft, Phys. Rev. Lett. 37 (1976) 8.
- 9) This was suggested to us by R. Barbieri.
- 10) A. De Rújula, Phys. Lett. 96B (1980) 279.
- 11) I.H. Dunbar and N. Dombey, University of Sussex preprint (1981).