Particles and Fields 2

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PREFACE

Particle physics seems to be entering a new period of consolidation. In 1977 when the first summer institute on particles and fields was held at the Banff Center, the standard model of the electro-weak interaction was a promising model more or less confirmed; today it seems quite well-confirmed. QCD was considered as probably the correct theory of strong interactions; today most theorists take it for granted. What seems to be lacking are computational tools and strenuous experimental testing; the major ideas seem to exist. Thus, this is a particularly auspicious time for a review of the status of theoretical and experimental particle physics and field theory.

The lectures collected in this volume were presented from August 16 to August 27, 1981 at the Banff Center in Banff, Canada. The unifying theme was gauge fields and the topics covered dealt with electro-weak interactions, Q.C.D., sub-quarks and unified theories.

The format of the Institute was as follows: thirteen lecture series of two to four hours each given by S. Brodsky, D. Bryman, M. Chen, S. Coleman, M. Creutz, H. Harari, J. Iliopoulos, C.H. Llewellyn-Smith, P. Lepage, D. Perkins and L. Susskind. In addition there were nine seminars (one hour each) given by G. Bodwin, G. Bunce, M. Chen, N.G. Deshpande, N. Dombey, L. Hall, K. Kang, C.H. Llewellyn-Smith, S.Y. Lo, H. Rykaczewski, and P.J. Steinhardt. There were further informal seminars (held at the Center) and informal discussions (held on the hiking trails). These are not included in these proceedings.

Also omitted is the banquet speech by Dr. D.E. Berghofer, Assistant Deputy Minister, Department of Advanced Education and Manpower, Government of Alberta. For copyright reasons, the lectures by S. Coleman and M. Creutz are also not included. Since the seminar by G. Bunce dealt with an experiment in the planning stages, it has also been omitted.

It is a tribute to the lecturers that in spite of the beautiful surroundings, the magnificent weather and the long lecture sessions, their lectures were always well attended. For this we would like to thank the lecturers as well as the International Advisory Committee for their assistance and advice. The members of this committee were:

J.D. Bjorken
S.D. Drell
H. Harari
G. 'tHooft
C.S. Lam
Y. Nambu
L. O'Raifeartaigh
A. Salam
R. Stora
M.K. Sundaresan
J.C. Taylor
S. Ting

The detailed work was handled by the four man local Organizing Committee which consisted of:

A.Z. Capri
A.N. Kamal
Y. Takahashi
H. Umezawa

all from the Physics Department of the University of Alberta.

We were also very ably assisted by Georgette Jolicoeur, the conference secretary and mother to lost graduate students. She, together with Mary Yiu, did an admirable job of typing the manuscript. Further we would like to thank all the other "volunteers" from the physics department at Alberta who helped with many of the little tasks that could have added up to an impossibly large task without them.

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The Winspear Foundation
The Theoretical Physics Institute of the University of Alberta

The Institute was organized under the auspices of the Theoretical Physics Division of the Canadian Association of Physics. CAP attempts to hold a summer school annually. The royalties from the sale of the proceedings of this and previous schools go toward support of future Summer Institutes.

Anton Z. Capri
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CONTENTS

Part II - Seminars

UNIFICATION

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I. THE U(1) X SU(2) MODEL OF WEAK AND ELECTROMAGNETIC INTERACTIONS

The weak and electromagnetic interactions are described by a gauge theory based on the group U(1) x SU(2). The fundamental Lagrangian contains the following three types of fields:

(i) Gauge fields. We have four vector bosons, one for each generator of the algebra. After the spontaneous breaking three of them become massive ($W^{\pm}$, Z^{o}) while the fourth, the photon, remains massless.

(ii) Fermion matter fields. The basic unit is a "family" consisting of fifteen two-component complex spinor fields organized in four left-handed doublets and seven right-handed singlets. The prototype is the electron family

$$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L \; ; \; \begin{pmatrix} u_i \\ d_i \end{pmatrix}_L \qquad i = \text{blue, white, red}$$

$$e_R \; ; \; u_{iR}, \; d_{iR}$$

With the discovery of charm the second family (the muon family) is complete, but in the third one, that of the tau-lepton, the upper quark (the top) still refuses to appear. A remarkable property of each family is that the sum of the electric charges of its members vanishes. This turns out to be necessary for the cancellation of triangle anomalies in the Ward identities of axial currents and hence, for the construction of a renormalizable theory. The total number of families is not restricted. In fact, we know of no good reason why any, beyond the first one, should exist. From cosmological arguments (the observed helium abundance in the universe) and assuming that we understand the process of nucleosynthesis, we can put limits on the number of massless (or nearly massless) neutrino species. This upper bound turns out to be surprisingly small: ~3-4. Either there are no more families, or the other neutrinos further up are massive (or the argument fails). From strong interactions we can put an upper bound of sixteen quark flavors before we lose asymptotic freedom, but this may well be a theoretical prejudice. Finally there are some limits on the masses of heavy fermions, but I know of no way to translate them into bounds on the number of families.

(iii) Higgs scalar fields. We would be very happy if we could live with only the first two kinds of fields, but in fact we need a third one, the scalar Higgs fields. Their non-zero vacuum expectation values break the gauge symmetry spontaneously thus providing masses to $W^{\pm}, Z^{o}$ as well as the fermions. In the standard model this is accomplished with a complex doublet of scalar fields. At the end one neutral spin zero boson survives as a physical particle. There are no severe restrictions on its mass, but we shall come back to this problem presently.

Our confidence in this model is based on the following qualitative as well as quantitative predictions among which, all those which can be tested with existing machines, have been verified experimentally.

(i) The neutral currents: In general, we would expect, for every flavor, a parameter that determines the strength of the neutral current relatively to the charged one and another to fix the ratio of the vector and axial parts. In the simplest model, in which the breaking comes through isodoublet scalars, they are all expressible in terms of a single one, the angle θ_W. This is brilliantly confirmed by the fact that the values of θ_W from neutrino reactions and those from polarized electron scattering, coincide.

(ii) The charmed particles were predicted to decay preferentially to strange particles in agreement with observation.

(iii) As we mentioned already, each family must be complete. Therefore the discovery of a new lepton (τ) was interpreted as the opening of a third family. Indeed the b quark has already been established and the discovery of t will confirm, or invalidate, this picture.

A personal estimation for the evolution of the experimental situation concerning the U(1) x SU(2) predictions is the following:

For the moment what we know for sure is only that we have a Fermi theory of weak interactions with a very particular neutral current. Assuming C.V.C., the magnitude of the radiative corrections to β- and μ-decays tells us that this picture is going to change before a scale of the order of 150 GeV.

The next step will be the discovery of the W's and Z. The $\bar{p}p$ collider seems to be on schedule, so I estimate the required time as 2-3 years. By then we shall know that Yukawa was right.

A detailed study of their properties will reveal the existence of the Yang-Mills self-couplings. Since it is a difficult experiment, even for L.E.P. I guess about twelve years. Only then shall we know whether Glashow was right.

This is still not good enough. In order to prove that we have a renormalizable theory we must discover the Higgs scalars or something equivalent. I can offer no guess on how long it is going to take. If the Higgs is light it may be found earlier than the previous step and then we shall prove that Weinberg and Salam were right.

What are the limits on the mass of the Higgs scalar? In the standard model with only one doublet we find $m_H^2 \sim \lambda G_F^{-1}$ where λ is the four-Higgs self-coupling constant and G_F is the Fermi constant. Therefore m_H cannot be arbitrarily large, otherwise the Higgs scalars become strongly interacting and perturbation theory breaks down. From this argument we obtain a limit

$$m_H < 100 \text{ GeV} - 1 \text{ TeV} \quad . \tag{1}$$

We can also obtain a lower limit. In the tree approximation the spontaneous symmetry breaking is obtained by imposing that the potential have a minimum away from the origin in field space. We should however verify that this is not changed when higher order corrections are included. This turns out to give a lower limit for m_H of the order

$$m_H > 10 \text{ GeV}. \tag{2}$$

Given these limits $10 \text{ GeV} < m_H < 1 \text{ TeV}$ the search for the Higgs scalar can take one of the following two directions: Direct search. The most caracteristic signature of a Higgs boson is its coupling properties to fermions. Again in the minimal standard model it is predicted to have Yukawa couplings of the form:

$$\frac{m_\psi}{v} \bar{\psi} \psi \phi \tag{3}$$

where v is its vacuum expectation value and m_ψ the mass of the fermion ψ. The interesting feature of this formula is that it predicts a Higgs scalar which is coupled preferentially to the heaviest fermion flavor available. It is a quite unique property which will be easy to verify if the beast is found. The second method to look for Higgs is an indirect search. Since its presence is necessary for renormalizability we expect to find some processes whose behaviour will be sensitive to m_H. However, becasue of peculiar cancellations among graphs, no dramatic effect can be found in any measurable quantity. For example, the predicted $e^+e^- \to W^+W^-$ cross section with longitudinally polarized W's varies by only 5% when the Higgs mass moves from 10 GeV to 1 TeV and this is still among the most sensitive ones. An important question is the following: What if no Higgs scalar is found at LEP, or even higher energies? One consequence is that new strong interactions will appear at a scale of ~1 TeV. A particular realization of this program is provided by the model of Technicolor, but I shall not discuss it in these lectures.

II. BEYOND U(1) x SU(2) x SU(3)

There are several reasons why the gauge theory based on U(1) x SU(2) x SU(3) cannot be considered as the final theory. The most important is that it is not a unified theory at all. Each group factor comes with its own coupling strength. Even worse is the presence of the U(1) factor because U(1) permits any number of coupling constants. In a non-abelian group the coupling constant is fixed by the gauge boson self-coupling and it must be the same, except for Clebsh-Gordan coefficients, for every matter multiplet. For U(1), however, this is not so. In other words, the present theory does not explain why electric charge seems to be quantized and we see no particles with charge equal to πe. For all those reasons, and others which I skip, several theorists try to go beyond the standard model. The hypothesis of grand unification states that U(1) x SU(2) x SU(3) is the remnant of a larger, semi-simple group G which is spontaneously broken at very high energies. The scheme looks like:

$$
\begin{array}{ccc}
G \xrightarrow[M]{\quad \dots \quad} & U(1) \times SU(2) \times SU(3) & \\
& \Big\downarrow \; m_W \sim 100 \text{ GeV} & \\
& U(1) \times SU(3) & \qquad (4)
\end{array}
$$

where the breaking of G may be a multistage one and M is one (or several) characteristic mass scale. There are several consequences of this general idea, but before going into them, let me explain in a simple way the underlying dynamics which describes the breaking (4).

The problem is the following: Since G is supposed to be semi-simple, it has only one coupling constant. Then how come that we observe three distinct interactions at present energies? In order to answer this question let me go with some detail into the renormalization program of the theory. For pedagogical purposes let me explain the case of $G = SU(5)$. The changes for any other group are straightforward. Furthermore, let me ignore the second breaking in (4). This means that I shall have all fermions massless and I shall include only Higgs scalars belonging to the adjoint representation of SU(5). I start from the Lagrangian:

$$\mathcal{L} = -\frac{1}{4}\mathrm{Tr}\, G_{\mu\nu}G^{\mu\nu} + \bar{\psi}\, i\not{D}\psi + \frac{1}{2}\,\mathrm{Tr}(D_\mu\Phi D^\mu\Phi)$$

$$+ \frac{1}{2}\, m^2\, \mathrm{Tr}(\Phi^2) - \frac{h_1}{4}\,[\mathrm{Tr}(\Phi^2)]^2 - \frac{h_2}{2}\,\mathrm{Tr}(\Phi^4)\ . \qquad (5)$$

The Lagrangian (5) contains four independent parameters, namely m^2, h_1, h_2 and g. After spontaneous symmetry breaking via the translation $\Phi \to \Phi + v/\sqrt{2}\,\lambda_{24}$, where λ_{24} is the matrix which commutes with the generators of U(1) x SU(2) x SU(3)

$$\lambda_{24} = \frac{2}{\sqrt{15}}\begin{pmatrix} 1 & & & & \\ & 1 & & & \\ & & 1 & & \\ & & & -3/2 & \\ & & & & -3/2 \end{pmatrix} \qquad (6)$$

we find the following mass spectrum:

(i) $M^2 = 5/3\, v^2 g^2$ for the vector gauge bosons $\Phi_9 \ldots \Phi_{20}$

(ii) $M_3^2 = 1/3\, v^2 h_2$ for the scalars $\Phi_1 \ldots\ldots \Phi_8$

(iii) $M_2^2 = 4/3\, v^2 h_2$ for the scalars $\Phi_{21} \ldots\ldots \Phi_{23}$

(iv) $M_1^2 = 2v^2(h_1 + 7/15\, h_2)$ for Φ_{24}

In terms of the parameters of (5) v^2 is given, in the tree approximation, by:

$$v^2 = m^2[h_1 + \frac{7}{15}\, h_2]^{-1} \qquad (7)$$

We choose to renormalize the broken theory by imposing the following seven renormalization conditions:

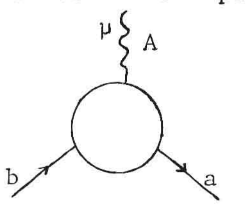
(i) Three conditions for the masses M^2, M_1^2 and M_2^2 which will be chosen physical, as the poles of the corresponding propagators.

(ii) Three wave function renormalization conditions for fermions, vector bosons and scalars respectively. For example the condition on the scalars reads:

$$\frac{d}{dp^2}\,\Gamma_A^{(2)}(p^2)\Big|_{p^2=-\mu^2} = 1 \qquad (8)$$

and similarly for fermions and gauge bosons. μ^2 is a substraction point. However, since SU(5) is broken, we must specify on which component of Φ we impose the condition. This is done with the index A which runs, in principle, from 1 to 24. The conservation of U(1) x SU(2) x SU(3) implies that we need only to distinguish four distinct cases (A = 1,...,8; 9,...,20; 21,...,23; 24). Once the condition is imposed on any one of the components all the other Green functions become finite and calculable. The same remarks apply to the vector bosons.

(iii) Finally, we must impose one coupling constant renormalization condition and we choose it to be the value of the fermion-vector boson three point function at $p_i^2 = -\mu^2$:



$$\Big|_{p_i^2 = -\mu^2} = i g_A \gamma^{\mu} T^A_{ab} \qquad (9)$$

where, again, the index A denotes the particular vector boson we have used.

I want to emphasize that, once these three kinds of renormalization conditions have been imposed, the perturbation theory is completely defined and all Green's functions are finite and calculable as formal power series in g_A. Furthermore, it is not possible to impose any further conditions. Notice, in particular, that there exists only one gauge coupling constant, as one should expect from a grand unified theory. Does this mean that we could compute a Q.C.D. process, like a deep inelastic structure function, as a power series in the weak interaction coupling constant? Formally, the answer is yes, but in practice this is not so. Formal perturbation theory guarantees that, if we choose in (9) A to denote one of the SU(2) gauge bosons, all other three point functions are finite and calculable. In particular, if $\Gamma_3^{(3)}$ is the three-point function with an SU(3) external gauge boson, we can write: (I shall use the following notation: For A = 1,...,8 I shall call $g_A = g_3$; A = 21,...,23 $g_A = g_2$; $g_{A=24} = g_1$):

$$\Gamma_3^{(3)}(p_i^2 = -\mu^2) = g_2 + A_1 g_2^3 + A_2 g_2^5 + \ldots \qquad (10)$$

where the A_n's are finite and calculable functions of μ^2 and the masses of the theory. However, it is easy to check that A_n is of the form:

$$A_n \sim \left(\ln \frac{M^2}{\mu^2}\right)^n \qquad (11)$$

which, for $M/\mu \sim 10^{14}$ gives $[65]^n$. In other words, although the series (10) is well defined, it is useless for practical computations.

The remedy to this difficulty is simple. We shall renormalize the same broken SU(5) theory in three different ways where the indices A in the conditions (8) and (9) denote the bosons of U(1), SU(2) or SU(3). This gives us three perturbation expansions in powers of g_1, g_2 or g_3, always of the same theory, but now each one is suited to particular processes. The values of the g_i's will be fixed by experiment, but we must always remember that we are talking about one and the same theory, so we write the analog of eq. (10):

$$g_i = F_{ij}(g_j, \frac{M^2}{\mu^2}, \alpha) \tag{12}$$

where with α we mean the ratios M_1^2/M^2 and M_2^2/M^2. In the limit of exact SU(5) symmetry, i.e. when $M^2/\mu^2 \to 0$, all coupling constants must be equal. This only happens at infinite energy

$$F_{ij}(g_j, 0, \alpha) = g_j \tag{13}$$

By taking $\mu^2 d/d\mu^2$ on both sides of (12) we obtain

$$\beta_i(F_{ij}, \lambda, \alpha) = [-\lambda \frac{\partial}{\partial\lambda} + \beta_j(g_j, \lambda, \alpha) \frac{\partial}{\partial g_j}]F_{ij} = 0 \tag{14}$$

where

$$\beta_k(g_k, \lambda, \alpha) = \mu^2 \frac{d}{d\mu^2} g_k \tag{15}$$

and $\lambda = M^2/\mu^2$.

The differential equation (14) with the boundary condition (13) is our basic equation. The β-functions are calculable at any given order of perturbation theory and they are of the form:

$$\beta_k(g_k, \lambda, \alpha) = b_k^o (\lambda, \alpha) g_k^3 + \ldots \tag{16}$$

Notice that the b coefficients, unlike those of F_{ij} itself, do not contain large logarithms. This can be easily understood since the β-functions, as defined by eqs. (15) and (9), possess well-defined limits both for $M^2/\mu^2 \to \infty$(when they become the β-functions of SU(3), SU(2) or U(1)) and $M^2/\mu^2 \to 0$ when they become all equal to that of SU(5). This is a consequence of the decoupling theorem. On the contrary, F_{ij} has no limit when $M^2/\mu^2 \to \infty$ with g_j kept fixed.

Equation (14) can be solved by the standard methods of characteristics. The solution expresses any coupling constant in terms of any other.

$$g_i = F_{ij} = \eta(g_j, \lambda, \alpha) \tag{17}$$

with η a given function. For example, using the one-loop β-functions of eq. (16) we find:

$$\frac{1}{g_i^2} = \frac{1}{g_j^2} + 2 \int_0^{\lambda} \frac{dx}{x} \, [b_i^o(x,\alpha) - b_j^o(x,\alpha)] \tag{18}$$

We must now use as input the experimentally measured effective strengths of strong, e.m. and weak interactions at moderate (say $p^2 \sim 10\text{-}100\ \mathrm{GeV}^2$) energies. Using the renormalization group for $\mu^2 = -p^2$, we write $g_i = \bar{g}_i(p^2)$ and $\lambda = -\mu^2/p^2$. The two independent equations given by the relations (17) $(i,j = 1,2,3)$ contain three unknown parameters, namely λ and the Higgs masses M_1^2/M^2 and M_2^2/M^2 denoted by α. However, it turns out that the dependence on α is very weak. If we ignore it for the moment, then the two equations can be used to determine λ and to predict the value of $\sin\theta_W$. A precise calculation must take into account the breaking of $U(1) \times SU(2)$ as well. In fact it turns out that the value of $\sin\theta_W$ is quite sensitive to this last breaking. The result, including the two-loop effects is (at $\mu = 3$ GeV):

$$\sin^2\theta_W = 0.216, \qquad M = 3.1\ 10^{14}\ \mathrm{GeV} \tag{19}$$

where the value of the Q.C.D. scale parameter $\Lambda \sim 0.2$ GeV has been used.

In Figure 1a we show the variation of the effective coupling

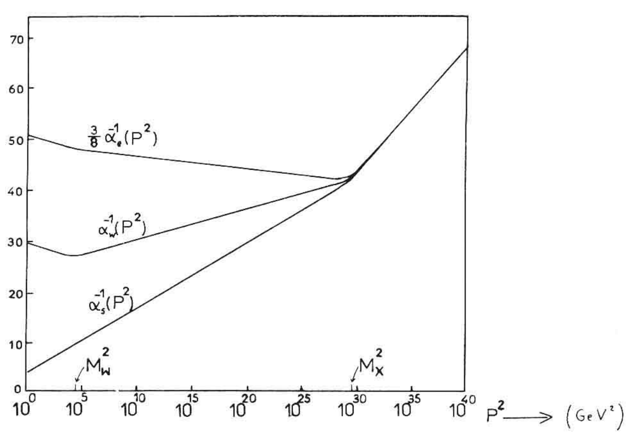


Fig. 1a. The variation of the effective coupling constants with the energy scale.