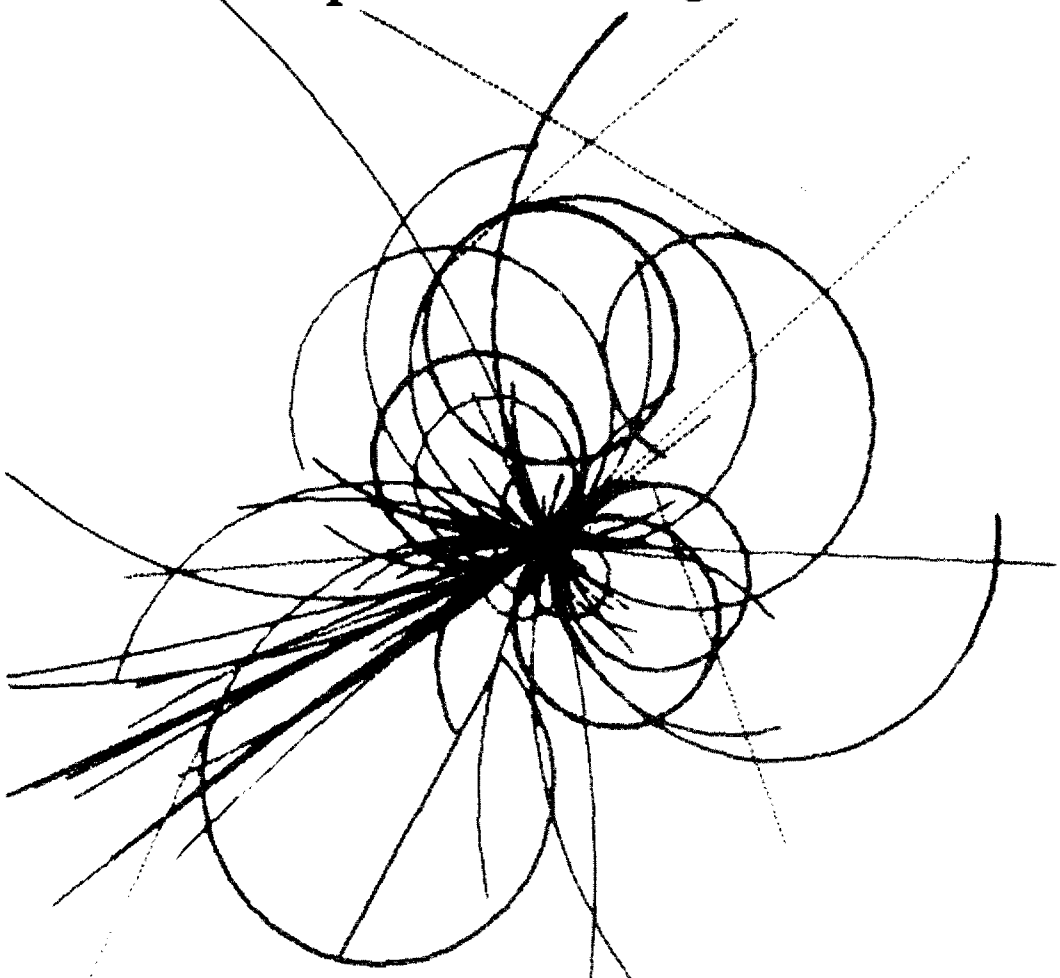
SSC-158

The Superconducting Super Collider



The Significance of the 1-TeV Scale

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Abstract

The Standard Model of elementary particle physics organizes our knowledge of the basic constituents and interactions, and achieves a wide-ranging synthesis of elementary phenomena, but leaves unanswered some fundamental questions. Important clues to a more complete understanding are to be found on the energy scale of 1 TeV, where the kinship between the weak and electromagnetic interactions will be illuminated. Exploration of the 1-TeV scale is an essential goal for the 1990s.

1 Introduction

In this third lecture of the 1987 Bernard Gregory Lectures, I wish to summarize the case for an experimental assault on the 1-TeV scale, a challenge that engages the imagination of accelerator designers the world around. I shall survey some of the reasons we believe a thorough study of collisions of the fundamental constituents at energies around 1 TeV is crucial to an understanding of electroweak symmetry breaking, and thus to the creation of a conceptual structure that goes beyond today's Standard Model of Elementary Particle Physics. My emphasis will be on motivations and goals, not on means, but I will comment briefly near the end of the lecture on some of the challenges of experimentation at multi-TeV hadron colliders.

The means of experimentation are of course highly important. The progress of our field can be chronicled in terms of the progress of accelerators and detectors, which is nourished by basic scientific discovery and advances in technology. A similar interplay between science and technology occurs in all fields of research. I have elaborated on the synergism between experiment, theory, and technology in my Gregory Lecture at Ecole Polytechnique.

In some cases, the parameters of new accelerators are sharply defined. CERN's current construction project is an apt example. Both phases of the electron-positron collider LEP respond to specific scientific demands: first, to produce a

large number of Z^0s in the formation reaction

$$e^+e^- \to Z^0 \,, \tag{1.1}$$

and later to surpass the threshold for the reaction

$$e^+e^- \to W^+W^- \,. \tag{1.2}$$

In other cases, the definition of parameters follows from more diffuse imperatives: to take a large step into unexplored territory, or to push the limits of an emerging technology. Though not characterized by a sharp threshold, the 1-TeV scale is an important landmark, and an essential goal for colliders of the next generation.

The picture[1] of the fundamental constituents of matter and the interactions among them that has emerged in recent years is one of great beauty and simplicity. All matter appears to be composed of *quarks* and *leptons*, which are pointlike, structureless, spin-$\frac{1}{2}$ particles. If we leave aside gravitation, which is a negligible perturbation at the energy scales usually considered, the interactions among these particles are of three types: weak, electromagnetic, and strong. All three of these interactions are described by gauge theories and are mediated by spin-1 gauge bosons. The quarks experience all three interactions; the leptons participate only in the weak and electromagnetic interactions. By the Standard Model we will understand two elements:

- The $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge theory of the strong, weak, and electromagnetic interactions; and
- Three generations of color-triplet quarks: *u, d; s, c; b, [t]*; and color-singlet leptons: ν_e, e; ν_μ, μ; ν_τ, τ.

The Standard Model has an appealing simplicity and an impressive generality. The picture at which we have arrived has a pleasing degree of coherence, and holds the promise of deeper understanding — in the form of a further unification of the interactions — still to come.

This is an accomplishment worthy of the pleasure we take in it, but if we have come impressively far in the past two decades, we still have quite far to go. The very success of the standard $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ model prompts new questions: Why does it work? Can it be complete? Where will it fail? As we shall see, the Standard Model itself hints that the frontier of our ignorance lies at $\sim$ 1 TeV for collisions among the fundamental constituents. In more general terms, the success of the Standard Model suggests that a significant step beyond present-day energies is needed, to see where our current understanding breaks down. A high-luminosity, multi-TeV proton-proton collider is the most technically assured and cost-effective instrument to take such a step, and to make

possible a timely and thorough exploration of the 1-TeV scale. The essential elements of supercollider physics are covered at greater length in *EHLQ*.[2] Many details are investigated in the proceedings of various workshops,[3–8] and in lecture notes from summer schools.[9–12]

2 What Landmarks Do We Expect?

I have already remarked in my introductory comments on the importance of the 1-TeV scale. In this section, I wish to review for the first time some of the arguments that lead to an identification of the 1-TeV scale as a key landmark. As we shall see again and again in different ways, our understanding of the spontaneous breaking of the electroweak gauge symmetry is incomplete. A more complete understanding can be obtained only with the aid of a thorough knowledge of what takes place on the 1-TeV scale.

Let us review the essential elements of the Weinberg-Salam $SU(2)_L \otimes U(1)_Y$ model of weak and electromagnetic interactions. To save writing, we shall speak of the model as it applies to a single generation of quarks and leptons. The generalization to several generations is well known.

We begin by specifying the fermions. The leptonic sector consists of a left-handed weak isospin doublet

$$\mathsf{L}_\ell = \begin{pmatrix} \nu_e \\ e \end{pmatrix}_L \tag{2.1}$$

with weak hypercharge $Y(\ell_L) = -1$, and a right-handed weak isospin singlet

$$\mathsf{R}_e \equiv e_R \tag{2.2}$$

with weak hypercharge $Y(e_R) = -2$. The hadronic sector is built upon a color triplet left-handed weak-isospin doublet

$$\mathsf{L}_q = \begin{pmatrix} u \\ d \end{pmatrix}_L \tag{2.3}$$

with weak hypercharge $Y(q_L) = 1/3$, and two color triplet right-handed weak-isospin singlets

$$\left.\begin{aligned} \mathsf{R}_u &\equiv u_R \\ \mathsf{R}_d &\equiv d_R \end{aligned}\right\} \tag{2.4}$$

with weak hypercharge $Y(u_R) = 4/3$ and $Y(d_R) = -2/3$.

The electroweak gauge group, $SU(2)_L \otimes U(1)_Y$, implies two sets of gauge fields: a weak isovector $\vec{b}_\mu$, with coupling constant g, and a weak isoscalar

$\mathcal{A}_\mu$, with coupling constant g'. Corresponding to these gauge fields are the field-strength tensors $\vec{F}_{\mu\nu}$ for the weak-isospin symmetry and $f_{\mu\nu}$ for the weak-hypercharge symmetry.

We may summarize the interactions by the Lagrangian

$$\mathcal{L} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{leptons}} \ , \tag{2.5}$$

with

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4} F^{\ell}_{\mu\nu} F^{\ell\mu\nu} - \frac{1}{4} f_{\mu\nu} f^{\mu\nu}, \tag{2.6}$$

and

$$\begin{aligned} \mathcal{L}_{\text{leptons}} &= \overline{\mathsf{R}}\, i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2}\mathcal{A}_\mu Y \right) \mathsf{R} \\ &+ \overline{\mathsf{L}}\, i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2}\mathcal{A}_\mu Y + i\frac{g}{2}\vec{\tau}\cdot\vec{b}_\mu \right) \mathsf{L}. \end{aligned} \tag{2.7}$$

To hide the electroweak symmetry, we introduce a complex doublet of scalar fields

$$\phi \equiv \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} \tag{2.8}$$

with weak hypercharge $Y_\phi = +1$. Add to the Lagrangian new terms for the interaction and propagation of the scalars,

$$\mathcal{L}_{\text{scalar}} = (\mathcal{D}^\mu \phi)^\dagger (\mathcal{D}_\mu \phi) - V(\phi^\dagger \phi), \tag{2.9}$$

where the gauge-covariant derivative is

$$\mathcal{D}_\mu = \partial_\mu + i\frac{g'}{2}\mathcal{A}_\mu Y + i\frac{g}{2}\vec{\tau}\cdot\vec{b}_\mu \ , \tag{2.10}$$

and the potential interaction has the form

$$V(\phi^\dagger\phi) = \mu^2(\phi^\dagger\phi) + |\lambda|\,(\phi^\dagger\phi)^2. \tag{2.11}$$

We are also free to add a Yukawa interaction between the scalar fields and the leptons:

$$\mathcal{L}_{\text{Yukawa}} = -G_e \left[\overline{\mathsf{R}}(\phi^\dagger \mathsf{L}) + (\overline{\mathsf{L}}\phi)\mathsf{R} \right]. \tag{2.12}$$

The electroweak symmetry is spontaneously broken if the parameter $\mu^2 < 0$. The minimum energy, or vacuum state, may then be chosen to correspond to the vacuum expectation value

$$\langle \phi \rangle_0 = \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix}, \tag{2.13}$$

where

$$v = \sqrt{-\mu^2/|\lambda|} = \left(G_F\sqrt{2}\right)^{-\frac{1}{2}} \approx 246 \text{ GeV} \tag{2.14}$$

is fixed by the low-energy phenomenology of charged current interactions.

The spontaneous symmetry breaking has several important consequences:

- Electromagnetism is mediated by a massless photon, coupled to the electric charge;
- The mediator of the charged-current weak interaction acquires a mass characterized by $M_W^2 = \pi\alpha/G_F\sqrt{2}\sin^2\theta_W$, where θ_W is the weak mixing angle;
- The mediator of the neutral current weak interaction acquires a mass characterized by $M_Z^2 = M_W^2/\cos^2\theta_W$;
- A massive neutral scalar particle, the Higgs boson, appears, but its mass is not predicted;
- Fermions can acquire mass.

Before reviewing the significance of the 1-TeV scale, it will be useful to recall why a Higgs boson, or its Doppelgänger, must exist. One path to the (theoretical!) discovery of the Higgs boson involves the role of the Higgs boson in the cancellation of high-energy divergences. An illuminating example is provided by the reaction

$$e^+e^- \to W^+W^-, \tag{2.15}$$

which is described in lowest order in the Weinberg-Salam theory by the four Feynman graphs in Figure 1. The leading divergence in the $J = 1$ amplitude of the neutrino-exchange diagram in Figure 1(a) is cancelled by the contributions of the direct-channel γ- and Z^0-exchange diagrams of Figs. 1(b) and (c). However, the $J = 0$ scattering amplitude, which exists in this case because the electrons are massive and may therefore be found in the "wrong" helicity state, grows as $s^{1/2}$ for the production of longitudinally polarized gauge bosons. The resulting divergence is precisely cancelled by the Higgs boson graph of Figure 1(d). If the Higgs boson did not exist, we should have to invent something very much like it. From the point of view of S-matrix theory, the Higgs-electron-electron coupling must be proportional to the electron mass, because "wrong helicity" amplitudes are always proportional to the fermion mass.

Let us summarize: Without spontaneous symmetry breaking in the Standard Model, there would be no Higgs boson, no longitudinal gauge bosons, and no

extreme divergence difficulties. (Nor would there be a viable low-energy phenomenology of the weak interactions.) The most severe divergences are eliminated by the gauge structure of the couplings among gauge bosons and leptons. A lesser, but still potentially fatal, divergence arises because the electron has acquired mass — because of the Higgs mechanism. Spontaneous symmetry breaking provides its own cure by supplying a Higgs boson to remove the last divergence. A similar interplay and compensation must exist in any satisfactory theory.

It is well known that the Standard Model does not give a precise prediction for the mass of the Higgs boson. We can, however, use arguments of self-consistency to place plausible lower and upper bounds on the mass of the Higgs particle in the minimal model. A lower bound is obtained by computing[13] the first quantum corrections to the classical potential (2.11). Requiring that $\langle\phi\rangle_0 \neq 0$ be an absolute minimum of the one-loop potential yields the condition

$$\begin{aligned} M_H^2 &> 3G_F\sqrt{2}(2M_W^4 + M_Z^4)/16\pi^2 \\ &\gtrsim 7\ \mathrm{GeV}/c^2\,. \end{aligned} \tag{2.16}$$

Unitarity arguments[14] lead to a conditional upper bound on the Higgs boson mass. It is straightforward to compute the amplitudes $\mathcal{M}$ for gauge boson scattering at high energies, and to make a partial-wave decomposition, according to

$$\mathcal{M}(s,t) = 16\pi \sum_J (2J+1) a_J(s) P_J(\cos\theta)\,. \tag{2.17}$$

Most channels "decouple," in the sense that partial-wave amplitudes are small at all energies (except very near the particle poles, or at exponentially large energies), for any value of the Higgs boson mass M_H. Four channels are interesting:

$$W_L^+W_L^-,\quad Z_L^0Z_L^0/\sqrt{2},\quad HH/\sqrt{2},\quad HZ_L^0\,, \tag{2.18}$$

where the subscript L denotes the longitudinal polarization states, and the factors of $\sqrt{2}$ account for identical particle statistics. For these, the s-wave amplitudes are all asymptotically constant (i.e., well behaved) and proportional to $G_FM_H^2$. In the high-energy limit,

$$\lim_{s\gg M_H^2}(a_0) \to \frac{-G_FM_H^2}{4\pi\sqrt{2}}\cdot\begin{bmatrix} 1 & 1/\sqrt{8} & 1/\sqrt{8} & 0 \\ 1/\sqrt{8} & 3/4 & 1/4 & 0 \\ 1/\sqrt{8} & 1/4 & 3/4 & 0 \\ 0 & 0 & 0 & 1/2 \end{bmatrix}. \tag{2.19}$$

Requiring that the largest eigenvalue respect the partial-wave unitarity condition $|a_0| \leq 1$ yields

$$M_H < \left(\frac{8\pi\sqrt{2}}{3G_F}\right)^{1/2} = 1\ \mathrm{TeV}/c^2 \tag{2.20}$$

as a condition for perturbative unitarity.

If the bound is respected, weak interactions remain weak at all energies, and perturbation theory is everywhere reliable. If the bound is violated, perturbation theory breaks down, and weak interactions among $W^{\pm}$, Z, and H become strong on the 1-TeV scale. This means that the features of strong interactions at GeV energies will come to characterize electroweak gauge boson interactions at TeV energies. We interpret this to mean that new phenomena are to be found in the electroweak interactions at energies not much larger than 1 TeV.

Does this analysis mean that the observation of a light Higgs boson would remove the motivation for exploring the 1-TeV scale? Decidedly not! The Standard Model is unnatural and gives us no means of understanding why a light Higgs boson should emerge. The problem is nicely illustrated by an analysis carried out by 't Hooft.[15]

Consider the Lagrangian $\mathcal{L}(\Lambda)$ as an effective field theory that describes physics at the shortest distances probed (characterized by an energy scale Λ) and at all longer distances in terms of the fields or degrees of freedom appropriate to that scale. At a higher energy scale Λ', the appropriate Lagrangian $\mathcal{L}(\Lambda')$ may involve different degrees of freedom. In this sense, any Lagrangian we encounter should be thought of as an effective Lagrangian describing physics in terms of the degrees of freedom characteristic of the highest energy scale probed by experiment. In spite of occasional assertions by some of our visionary colleagues, we can never be certain that we have encountered all the fundamental fields that are to be discovered, up to the highest energies.

What properties must an effective Lagrangian display in order that it can consistently represent the low-energy effective interactions of some unknown dynamics acting at a higher energy scale? We say that the Lagrangian $\mathcal{L}(\Lambda)$ is *natural* if every small parameter ξ (in units of the requisite power of Λ) is associated with an approximate symmetry of $\mathcal{L}(\Lambda)$ that becomes an exact symmetry in the limit $\xi \to 0$. This is to say that dynamical accidents are unnatural and unsatisfying.

This definition has two important virtues:

- To determine whether a theory is natural at a scale Λ requires no knowledge of physics at scales above Λ.
- A Lagrangian unnatural on a scale Λ is unnatural on all higher scales $\Lambda' > \Lambda$.

For the Weinberg-Salam Lagrangian, the only possible additional symmetry that could allow for a naturally small scalar mass is invariance under a shift in the scalar field ϕ, which would call for $|\lambda|$, $\mu \to 0$. Such a symmetry is broken

by gauge interactions and scalar self-interactions. The theory is natural only if

$$\frac{M_H^2}{\Lambda^2} \gtrsim O(|\lambda|),\ O(\alpha)\ . \tag{2.21}$$

Therefore, since $M_H^2 \approx \lambda v^2$, the theory is natural only for scales

$$\Lambda \lesssim O(v) \approx 246 \text{ GeV} \sim M_W\ . \tag{2.22}$$

Consequently we conclude that values $M_H \ll M_W$ are unnatural.

We shall next give a more operational discussion of the naturalness problem. For now, let us note that two strategies for resolving the unnaturalness of the electroweak theory suggest themselves:

- Eliminate the scalars as fundamental degrees of freedom in $\mathcal{L}$ for $\Lambda \gg G_F^{-\frac{1}{2}}$, as in theories of technicolor or compositeness.
- Associate an approximate symmetry with light scalars, as in supersymmetric theories.

Both alternatives require new physics at or below the 1-TeV scale.

3 Why there must be new physics on the 1-TeV scale

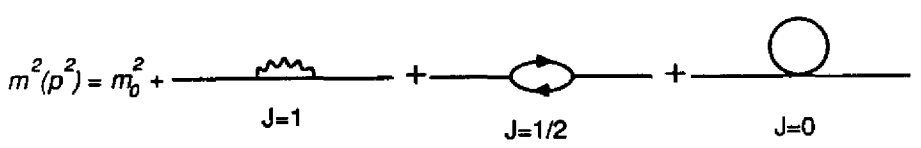
The Standard Model is incomplete[16]; it does not explain how the scale of electroweak symmetry breaking is maintained in the presence of quantum corrections. The problem of the scalar sector can be summarized neatly as follows.[17] The Higgs potential of the $SU(2)_L \otimes U(1)_Y$ electroweak theory is

$$V(\phi^\dagger\phi) = \mu^2(\phi^\dagger\phi) + |\lambda|\,(\phi^\dagger\phi)^2\ . \tag{3.1}$$

With μ_0^2 chosen to be less than zero, the electroweak symmetry is spontaneously broken down to the $U(1)$ of electromagnetism, as the scalar field acquires a vacuum expectation value fixed by the low-energy phenomenology,

$$\langle\phi\rangle_0 = \sqrt{-\mu^2/2|\lambda|} \equiv (G_F\sqrt{8})^{-1/2} \approx 175 \text{ GeV}\ . \tag{3.2}$$

Beyond the classical approximation, scalar mass parameters receive quantum corrections involving loops containing particles of spins $J = 1, 1/2$, and 0:



(3.3)

The loop integrals are potentially divergent. Symbolically, we may summarize the content of Eq. (3.3) as

$$m^2(p^2) = m^2(\Lambda^2) + Cg^2 \int_{p^2}^{\Lambda^2} dk^2 + \cdots , \tag{3.4}$$

where Λ defines a reference scale at which the value of m^2 is known, g is the coupling constant of the theory, and C is a constant of proportionality, calculable in any particular theory. Instead of dealing with the relationship between observables and parameters of the Lagrangian, we choose to describe the variation of an observable with the momentum scale. In order for the mass shifts induced by radiative corrections to remain under control (i.e., not to greatly exceed the value measured on the laboratory scale), either

- Λ must be small, so the range of integration is not enormous, or
- new physics must intervene to cut off the integral.

In the standard $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ model, the natural reference scale is the Planck mass,

$$\Lambda \sim M_{\text{Planck}} \approx 10^{19} \text{ GeV} . \tag{3.5}$$

In a unified theory of the strong, weak, and electromagnetic interactions, the natural scale is the unification scale,

$$\Lambda \sim M_U \approx 10^{15} \text{ GeV} . \tag{3.6}$$

Both estimates are very large compared to the scale of electroweak symmetry breaking (3.2). We are therefore assured that new physics must intervene at an energy of approximately 1 TeV, in order that the shifts in m^2 not be much larger than (3.2).

Only a few distinct classes of scenarios for controlling the contribution of the integral in (3.4) can be envisaged. The supersymmetric solution[18] is especially elegant. Exploiting the fact that fermion loops contribute with an overall minus sign (because of Fermi statistics), supersymmetry balances the contributions of fermion and boson loops. In the limit of unbroken supersymmetry, in which the masses of bosons are degenerate with those of their fermion counterparts, the cancellation is exact:

$$\sum_{i=\substack{\text{fermions}\\+\text{bosons}}} C_i \int dk^2 = 0 . \tag{3.7}$$

If the supersymmetry is broken (as it must be in our world), the contribution of the integrals may still be acceptably small if the fermion-boson mass splittings ΔM are not too large. The condition that $g^2 \Delta M^2$ be "small enough" leads to the requirement that superpartner masses be less than about 1 TeV/c^2.

A second solution to the problem of the enormous range of integration in (4.4) is offered by theories of dynamical symmetry breaking such as technicolor.[19] In the technicolor scenario, the Higgs boson is composite, and new physics arises on the scale of its binding, $\Lambda_{TC} \simeq O(1 \text{ TeV})$. Thus the effective range of integration is cut off, and mass shifts are under control.

A third possibility, which is appealingly economical but entails the sacrifice of perturbation theory for the electroweak interactions, is that of a strongly interacting gauge sector.[20] This would give rise to WW resonances, multiple production of gauge bosons, and other new phenomena.

Nature may choose any (or none) of these human inventions, but we are driven unavoidably to the conclusion that some new physics must occur on the 1-TeV scale.

4 Technicolor

No direct experimental evidence compels the modification or extension of the Standard Model. The motivations for going beyond the Standard Model, or for attempting to "complete" it, are based upon aesthetic principles of theoretical simplicity and elegance, or demands for internal consistency. Having reviewed some of the arguments for elaborating upon the Standard Model, I now consider one example of several possible extensions: the technicolor scheme of dynamical symmetry breaking. I select this in part because the other leading candidate, supersymmetry, is so well known, and in part because I find its claim on our attention very powerful.

We are not looking for a replacement of the Standard Model, for we expect that the Standard Model will remain as the low-energy limit of a more complete theory, much as the four-fermion description of the charged current weak interaction emerges as the low-energy limit of the Weinberg-Salam model.

4.1 The Idea of Technicolor

The dynamical symmetry-breaking approach, of which technicolor theories are exemplars, is modeled upon our understanding of another manifestation of spontaneous symmetry breaking in nature, the superconducting phase transition. The macroscopic order parameter of the Ginzburg-Landau phenomenology[21] corresponds to the wave function of superconducting charge carriers. It acquires a nonzero vacuum expectation value in the superconducting state. The microscopic Bardeen-Cooper-Schrieffer theory[22] identifies the dynamical origin of the order parameter with the formation of bound states of elementary fermions, the Cooper pairs of electrons. The basic idea of the technicolor mechanism

is to replace the elementary Higgs boson of the Standard Model by a fermion-antifermion bound state. By analogy with the superconducting phase transition, the dynamics of the fundamental technicolor gauge interactions among technifermions generate scalar bound states, and these play the role of the Higgs fields.

In the case of superconductivity, the elementary fermions (electrons) and the gauge interactions (QED) needed to generate the scalar bound states are already present in the theory. Could we achieve a scheme of similar economy for the electroweak symmetry-breaking transition?

Consider an $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ theory of massless up and down quarks. Because the strong interaction is strong, and the electroweak interaction is feeble, we may consider the $SU(2)_L \otimes U(1)_Y$ interaction as a perturbation. For vanishing quark masses, QCD has an exact $SU(2)_L \otimes SU(2)_R$ chiral symmetry. At an energy scale $\sim \Lambda_{QCD}$, the strong interactions become strong, fermion condensates appear, and the chiral symmetry is spontaneously broken

$$SU(2)_L \otimes SU(2)_R \rightarrow SU(2)_V \tag{4.1}$$

to the familiar flavor symmetry. Three Goldstone bosons appear, one for each broken generator of the original chiral invariance. These were identified by Nambu[23] as three massless pions.

The broken generators are three axial currents whose couplings to pions are measured by the pion decay constant f_π. When we turn on the $SU(2)_L \otimes U(1)_Y$ electroweak interaction, the electroweak gauge bosons couple to the axial currents and acquire masses of order $\sim g f_\pi$. The massless pions thus disappear from the physical spectrum, having become the longitudinal components of the weak gauge bosons. This achieves much of what we desire. Unfortunately, the mass acquired by the intermediate bosons is far smaller than required for a successful low-energy phenomenology; it is only[24]

$$M_W \sim 30 \text{ MeV}/c^2 \,. \tag{4.2}$$

4.2 A MINIMAL MODEL

The simplest transcription of these ideas to the electroweak sector is the minimal technicolor model of Weinberg[25] and Susskind.[26] The technicolor gauge group is taken to be $SU(N)_{TC}$ (usually $SU(4)_{TC}$), so the gauge interactions of the theory are generated by

$$SU(4)_{TC} \otimes SU(3)_c \otimes SU(2)_L \otimes U(1)_Y \,. \tag{4.3}$$

The technifermions are a chiral doublet of massless color singlets

$$\begin{pmatrix} U \\ D \end{pmatrix}_L \quad U_R,\ D_R\ . \tag{4.4}$$

With the electric charge assignments $Q(U) = \frac{1}{2}$ and $Q(D) = -\frac{1}{2}$, the theory is free of electroweak anomalies. The ordinary fermions are all technicolor singlets.

In analogy with our discussion of chiral symmetry breaking in QCD, we assume that the chiral TC symmetry is broken,

$$SU(2)_L \otimes SU(2)_R \otimes U(1)_V \to SU(2)_V \otimes U(1)_V\ . \tag{4.5}$$

Three would-be Goldstone bosons emerge. These are the technipions

$$\pi_T^+,\quad \pi_T^0,\quad \pi_T^-, \tag{4.6}$$

for which we are free to *choose* the technipion decay constant as

$$F_\pi = \left(G_F\sqrt{2}\right)^{-1/2} = 247\ \text{GeV}\ . \tag{4.7}$$

When the electroweak interactions are turned on, the technipions become the longitudinal components of the intermediate bosons, which acquire masses

$$\begin{aligned} M_W^2 &= g^2F_\pi^2/4 &&= \frac{\pi\alpha}{G_F\sqrt{2}\sin^2\theta_W} \\ M_Z^2 &= (g^2+g'^2)\,F_\pi^2/4 &&= M_W^2/\cos^2\theta_W \end{aligned} \tag{4.8}$$

that have the canonical Standard Model values, thanks to our choice (4.7) of the technipion decay constant.

Working by analogy with QCD, we may guess the spectrum of other $F\overline{F}$ bound states as follows:

$$\left.\begin{array}{cc} 1^{--}\ \text{technirhos} & \rho_T^+,\ \rho_T^0,\ \rho_T^- \\ 1^{--}\ \text{techniomega} & \omega_T \\ 0^{-+}\ \text{technieta} & \eta_T \\ 0^{++}\ \text{technisigma} & \sigma_T \end{array}\right\}, \tag{4.9}$$

all with masses on the order of the technicolor scale $\Lambda_{TC} \sim O(1\ \text{TeV}/c^2)$, since they do not originate as Goldstone bosons. The dominant decay of the technirho will be

$$\rho_T \to \pi_T\pi_T\ , \tag{4.10}$$

i.e., into pairs of longitudinally polarized gauge bosons. Standard estimates lead to

$$\begin{aligned} M(\rho_T) &\approx 1.77\ \text{TeV}/c^2 \\ \Gamma(\rho_T) &\approx 325\ \text{GeV}\ . \end{aligned} \tag{4.11}$$

4.3 Extended Technicolor

Technicolor shows how the generation of intermediate boson masses could arise without fundamental scalars or unnatural adjustments of parameters. It thus provides an elegant solution to the naturalness problem of the Standard Model. However, it has a major deficiency: it offers no explanation for the origin of quark and lepton masses, because no Yukawa couplings are generated between Higgs fields and quarks or leptons.

A possible approach to the problem of quark and lepton masses is suggested by "extended technicolor" models. We imagine that the technicolor gauge group is embedded in a larger extended technicolor gauge group,

$$G_{TC} \subset G_{ETC} \ , \tag{4.12}$$

which couples quarks and leptons to the technifermions. If the ETC symmetry is spontaneously broken down to the TC symmetry

$$G_{ETC} \to G_{TC} \ , \tag{4.13}$$

at a scale

$$\Lambda_{ETC} \sim 30 \ - \ 300 \text{ TeV} \ , \tag{4.14}$$

then the quarks and leptons may acquire masses

$$m \sim \Lambda_{TC}^3 / \Lambda_{ETC}^2 \ . \tag{4.15}$$

The outlines of this strategy are given in Refs. 27 and 28, but no "standard" ETC model has been constructed.

As a representative of the ETC strategy we may consider a model due to Farhi and Susskind.[29] Their model is built on new fundamental constituents, the techniquarks

$$\begin{pmatrix} U \\ D \end{pmatrix}_L \quad U_R, \ D_R \ , \tag{4.16}$$

which are analogs of the ordinary quarks, and the technileptons

$$\begin{pmatrix} N \\ E \end{pmatrix}_L \quad N_R, \ E_R \ , \tag{4.17}$$

which are analogs of the ordinary leptons. These technifermions are bound by the $SU(N)_{TC}$ gauge interaction, which is assumed to become strong at $\Lambda_{TC} \sim$ 1 TeV. Among the $F\overline{F}$ bound states are eight color-singlet, technicolor-singlet

pseudoscalar states [labeled by (I, I_3)]

$$\left.\begin{array}{ll} \pi_T^+ & (1,1) \\ \pi_T^0 & (1,0) \\ \pi_T^- & (1,-1) \end{array}\right\} \text{become longitudinal } W^\pm,\ Z^0$$

$$\left.\begin{array}{ll} P^+ & (1,1) \\ P^0 & (1,0) \\ P^- & (1,-1) \\ P^{0\prime} & (0,0) \end{array}\right\} \text{pseudo-Goldstone bosons} \qquad , \tag{4.18}$$

$$\eta_T' \quad (0,0) \qquad \text{techniflavor singlet}$$

plus the corresponding technivector mesons. Like the η' of QCD, the η_T' couples to an anomalous current, so it is expected to acquire a mass on the order of several hundred GeV/c^2. The pseudo-Goldstone bosons are massless in the absence of electroweak and ETC interactions.

The possibilities for study of the light particles implied in such a model have been examined recently by Eichten, Hinchliffe, Lane, and myself.[30] Some consequences of the extended technicolor interaction are examined in detail. In the absence of extended technicolor interactions, the neutral technipions P^0 and $P^{0\prime}$ remain massless, while the charged technipions P^+ and P^- acquire electroweak masses of a few GeV/c^2. When ETC interactions are included, the technipion masses have been estimated as[28,31]

$$\begin{aligned} 8 \ \text{GeV}/c^2 < \quad & M(P^\pm) \quad < 40 \ \text{GeV}/c^2 \,, \\ 2 \ \text{GeV}/c^2 < \quad & M(P^0, P^{0\prime}) \quad < 40 \ \text{GeV}/c^2 \,. \end{aligned} \tag{4.19}$$

If, as expected in the simplest models of Higgs bosons, the couplings of pseudoscalars into fermion pairs are proportional to fermion mass, the dominant decay modes will be

$$\begin{aligned} P^+ &\rightarrow t\bar{b}, c\bar{b}, c\bar{s}, \tau^+\nu \\ P^0 &\rightarrow b\bar{b}, c\bar{c}, \tau^+\tau^- \\ P^{0\prime} &\rightarrow b\bar{b}, c\bar{c}, \tau^+\tau^-;\ gg \,. \end{aligned} \tag{4.20}$$

Despite the possible similarities between Higgs bosons and technipions, there are important distinguishing characteristics. First, in the Standard Model, there is a direct HZZ coupling in the Lagrangian, whereas in the Farhi-Susskind model the P^0ZZ coupling is induced. As a consequence, we would expect the decay of a virtual $Z^* \rightarrow ZH$ to be about four orders of magnitude stronger than that of

$Z^* \to ZP^0$. If a Higgs-like entity is seen in the reaction

$$e^+e^- \to Z^0 \;\to\; Z^*H \quad \hookrightarrow \ell^+\ell^- \,, \tag{4.21}$$

then it is the Higgs, and technicolor is ruled out. Second, in a multi-Higgs model, the decay $Z^0 \to H^0H^{0\prime}$ is allowed (although the rate depends on details, such as mixing angles). In contrast, the decay $Z^0 \to P^0P^{0\prime}$ is inhibited; we therefore expect

$$\Gamma(Z^0 \to P^0P^{0\prime}) \ll \Gamma(Z^0 \to H^0H^{0\prime}) \,. \tag{4.22}$$

A clear presentation of the differences between Higgs bosons and technipions is given in Ref. 32.

4.4 Appraisal

Technicolor represents one of only a few known approaches to the problem of electroweak symmetry breaking. If the technicolor hypothesis correctly describes the breakdown of the electroweak gauge symmetry, there will be a number of spinless technipions with masses below the technicolor scale of about 1 TeV. Some of these, the color singlet, technicolor singlet particles, should be quite light (with masses $\lesssim 40$ GeV/c^2) and could be studied using the current generation of e^+e^- and $\bar{p}p$ colliders. Similar light scalars arise in multiple Higgs models and in supersymmetry. The colored particles are probably inacessible to experiment before a supercollider comes into operation, as are technivector mesons. Full exploitation of the scientific opportunities requires the efficient identification and measurement of heavy quark flavors, and the ability to identify intermediate bosons in complex events.

I am frequently asked what effect high–critical-temperature oxide superconductors will have on elementary particle physics. The questioner usually has in mind potential applications of a miraculous new substance to the technology of particle accelerators. This is an interesting area for speculation and dreaming, but I submit that it misses the truly revolutionary opportunities presented by the phenomenon of 1-2-3 superconductors.

The common thread of the progress in elementary particle physics over the past twenty-five years has been the shameless exploitation of ideas appropriated from condensed matter physics. The Ginzburg-Landau theory is none other than the nonrelativistic limit of the Abelian Higgs model, which opened the way to our understanding of spontaneous breaking of gauge symmetries. We have just seen that Technicolor theories draw inspiration from the BCS theory.

In this light, the implications of oxide superconductors for particle physics should be obvious to any ambitious young theorist. What we must do is identify

the correct theory of high-T_c superconductors ... and steal it!

5 EXPERIMENTAL ENVIRONMENT OF HADRON SUPERCOLLIDERS

What will experimentation be like at a multi-TeV proton-proton collider? Specific analysis of signals and backgrounds is quite fruitful, and has been the object of many of the studies carried out over the past four years.[2–8] However, we must also be aware of the general environment in which detectors must function and events must be selected and recorded.

The basic parameters of the Superconducting Super Collider are set out in the SSC Conceptual Design Report,[33] a non–site-specific conception of a 20⊕20 TeV proton-proton collider 83 km in circumference. The design calls for two clusters of interaction regions incorporating both physics experimental areas and major supporting equipment, a configuration that seems advantageous from the points of view of operating efficiency, economics, sociology, and accelerator physics. At the design luminosity of 10^{33} cm^{-2}sec^{-1}, interactions will occur at the rate of

$$0.016 \cdot (\sigma_t/1 \text{ mb}) \text{ interactions/crossing} . \tag{5.1}$$

The length of each bunch of protons is 6.0–7.3 cm, and adjacent bunches are separated by 5.1 m. A sketch of the layout proposed for the SSC is shown in Figure 2.

We expect the total cross section at 40 TeV to lie in the range

$$100 \text{ mb} \lesssim \sigma_t \lesssim 200 \text{ mb} \quad , \tag{5.2}$$

so that the event rate at the design luminosity may range up to 2×10^8 per second. A "best guess" for the total cross section, based on fits[34] to data up to $S\overline{p}pS$ energies, is $\sigma_t = 138$ mb .

A good way to gain respect for the conditions that will prevail at the SSC is to examine the trigger rate for events with transverse energy E_T greater than some threshold E_T^{min}. This is shown in Figure 3 for the nominal operating conditions of the SSC: $\sqrt{s} = 40$ TeV and $\mathcal{L} = 10^{33}$ cm^{-2}sec^{-1}, as well as at 10 and 100 TeV. At 40 TeV, a "high-E_T" trigger with threshold set at 2 TeV will count at 1 Hz from two-jet QCD events. This is of interest in planning triggers that will efficiently select interesting events from the approximately 140 million interactions that will take place each second in an SSC interaction region.

Particle multiplicities will also be large. In QCD, we can estimate[35] the multiplicity of partons in a gluon jet as

$$\langle n(Q) \rangle_g \propto (\log Q/\Lambda)^{-c} \exp\left(\sqrt{\frac{12 \log Q/\Lambda}{\pi b_0}}\right) , \tag{5.3}$$

where $Q \sim p_\perp$ measures the virtualness of the jet, and

$$b_0 = (33 - 2N_f)/12\pi \; ; \quad \text{and } c = (11 + 22N_f/27)/8\pi b_0 \; , \tag{5.4}$$

where N_f is the number of active quark flavors. This result depends critically on taking account of quantum mechanical interference; in a purely probabilistic ("branching") approach, the exponent is larger by a factor of $\sqrt{2}$.

The multiplicity difference between quark jets and gluon jets is also calculable in perturbative QCD. The result[36] is

$$\frac{\langle n(Q)\rangle_g}{\langle n(Q)\rangle_q} = \frac{9}{4} \cdot (1 - 0.27\sqrt{\alpha_s} - 0.07\alpha_s) \; . \tag{5.5}$$

The multiplicities expected in a 1-TeV jet are impressively large, on the order of 25–50.

Finally, the fragmentation of gluon jets into heavy quark pairs is expected to be reliably calculated in perturbation theory.[37] Roughly speaking, a 1-TeV gluon will yield 0.5 pair of c-quarks, 0.25 pair of b-quarks, and 0.05 t-quarks.

6 Concluding Remarks

In this lecture, I have reviewed the case for exploration of the 1-TeV scale. The description of the strong, weak, and electromagnetic interactions of the fundamental constituents — the quarks and leptons — in terms of gauge theories based on the symmetry group $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ is aesthetically appealing and has had many experimental successes. The unity and predictive power already achieved, and the promise of a more complete unification of the fundamental interactions, make it imperative to examine the foundations of the current paradigm and to take seriously its shortcomings as hints for improvements.

The incompleteness of our theoretical description is manifested by our ignorance of the dynamical mechanism that underlies the spontaneous breaking of electroweak gauge symmetry, by the multitude of seemingly arbitrary parameters required to specify the Standard Model, by the puzzling repetition of quark and lepton generations, and by many other questions.

In this brief survey, it has been possible only to scratch the surface of the physics opportunities presented by a high-energy, high-luminosity hadron collider. The scope of physics issues to be addressed ranges from detailed study of known particles, such as the intermediate bosons, to the search for high-mass exotica. The comprehensive studies of physics possibilities carried out over the past four years have shown convincingly that important clues are to be found on

the scale of 1 TeV, and that a high-luminosity, multi-TeV hadron supercollider will supply the means to reveal them.

With respect to supercollider experimentation, there are a few detector issues that I like to raise at every opportunity.

- The utility of high-efficiency W and Z detectors. The discovery physics we have considered in assessing the physics prospects of the SSC can all be done by relying upon the leptonic decays of the gauge bosons, but we can move to a deeper level of experimentation by learning to use the nonleptonic decays as well.

- The UA1 experiment has already indicated the value of "hermetic" detectors, which can capture and measure all the visible energy emitted in the central region. For a general-purpose SSC detector, it is of interest to require hermeticity for rapidities $|y| < 3$.

- Examples from technicolor and the Higgs sector of the Standard Model indicate that good-efficiency $\tau, b, \ldots$ tags will be of considerable value in enhancing signals over background. Full utilization of the heavy-flavor tag requires measuring the four-momenta of the short-lived particles as well.

- How to reduce the interaction rate of $\sim 10^8$ Hz to the $O(1\text{ Hz})$ rate at which complex events can be written on storage media (magnetic tapes, optical discs)?

- Bringing remote local intelligence into the detector components themselves requires the implementation of radiation-hardened electronics, especially near the beam directions.

We have recognized the significance of the 1-TeV scale for a decade. Through the development of superconducting magnets, and thanks to the experience gained in operating high-energy $\overline{p}p$ colliders at CERN and Fermilab and the evolution of detector architecture from Mark I at SPEAR up through the upgraded UA1 and UA2 at CERN and CDF at Fermilab, we now have the technical means in hand to begin our assault on this frontier of our ignorance.

We are faced with great opportunities!

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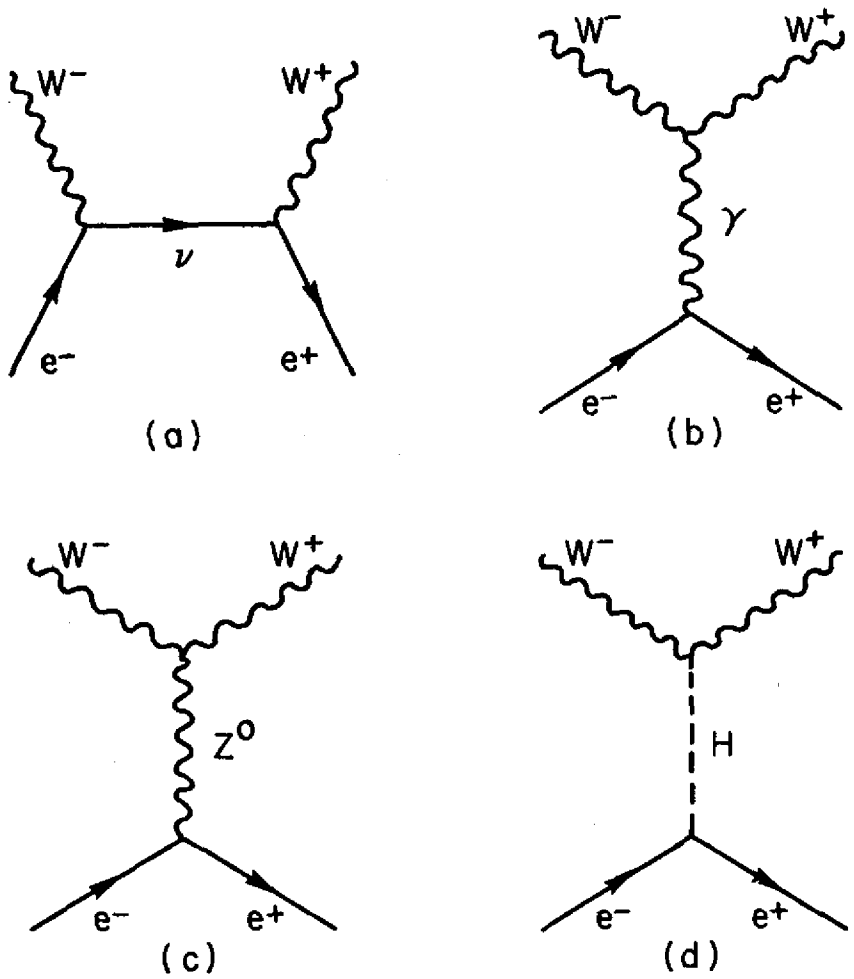

Figure 1: Lowest-order contributions to the reaction $e^+e^- \to W^+W^-$ in the Standard Model.

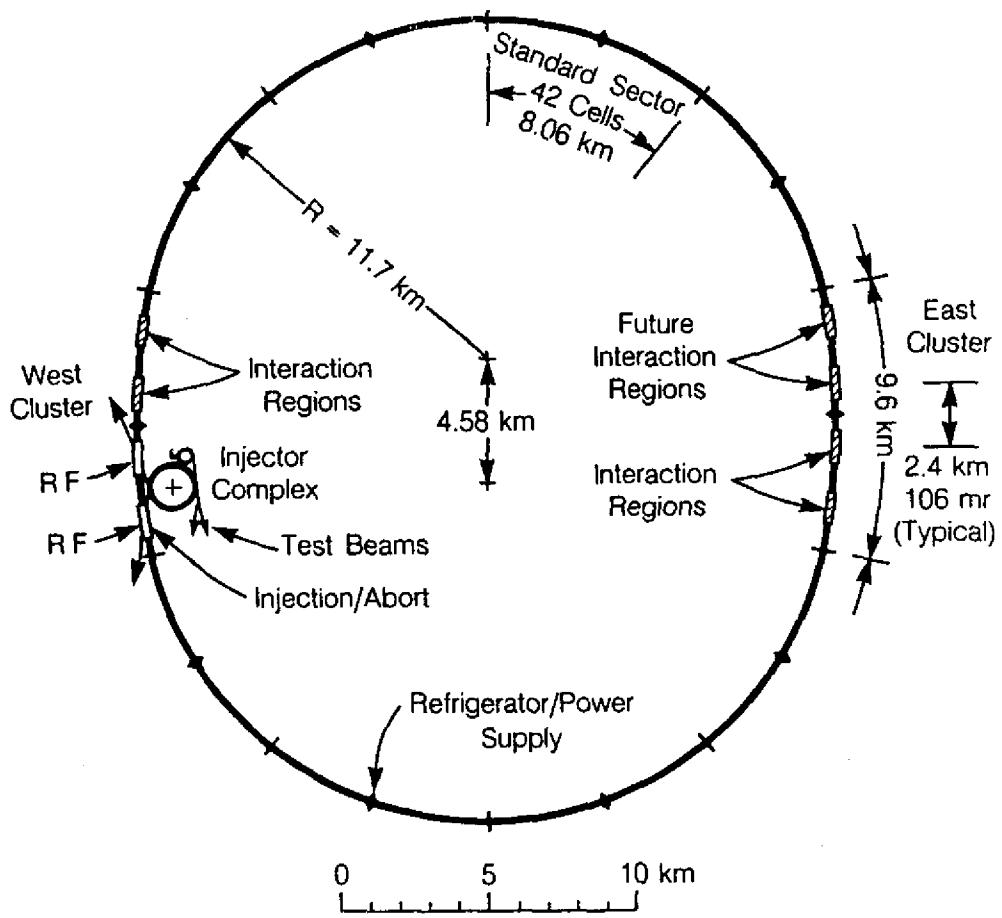

Figure 2: SSC collider ring layout. East and west clusters are joined by arcs of 11.7 km radius. The east cluster consists of four interaction regions separated by 2.4 km. The west cluster has two interaction regions and two utility straight sections (open rectangles) for injection and abort and for acceleration (RF). The cascade of synchrotrons that forms the injector is inside the main ring at the utility straight sections. There are 10 refrigeration and power units around the ring (black diamonds).

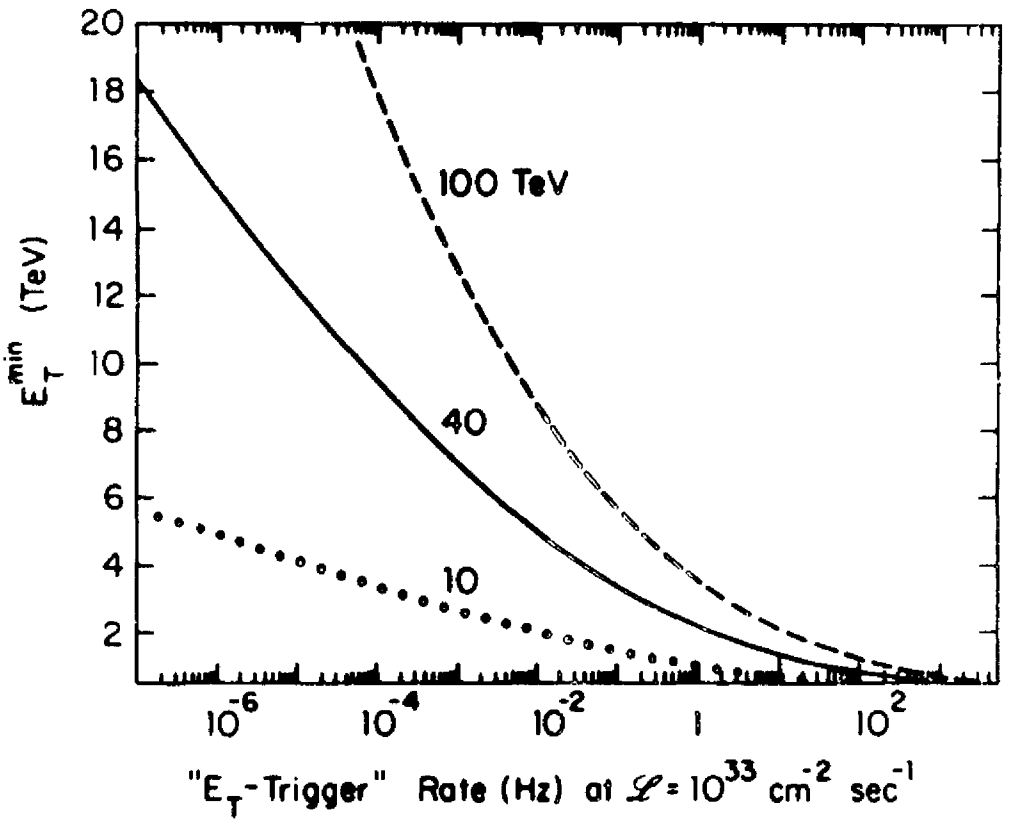

Figure 3: Counting rate for an E_T-trigger in pp collisions at an instantaneous luminosity of $\mathcal{L} = 10^{33}$ cm^{-2}sec^{-1} (after *EHLQ*). The threshold is defined for transverse energy deposited in the central region of rapidity, defined by $|y_i| < 2.5$ for jets 1 and 2.