Conference Summary

Steven Weinberg

Theory Group
Department of Physics
University of Texas
Austin, Texas 78712

Talk presented at the XXVI International Conference on High Energy Physics, Dallas, Texas, August, 12, 1992.

*Research supported in part by the Robert A. Welch Foundation and NSF Grant PHY 9009850.
Roy Schwitters has thanked the staff of this conference for their good work. As this is the last talk, I would like to add my own thanks to the staff, and also on behalf of all the participants to thank Roy himself and Vic Teplitz and the other physicists at the Super Collider and at SMU and the other Dallas area universities for their part in making this such an exceptionally well run conference.

This is actually the second time that I have been called on to give the summary talk at one of the Rochester Conferences. The previous time was at Berkeley in 1986, the latest time the conference was held in the United States. You might think that after that experience I would know better than to take on such a difficult task, but in fact nothing could be further from the truth. This is actually an extremely easy job. Everyone knows that it is impossible to review every topic that has been discussed at such a conference, so no one expects it. In fact I found last time at Berkeley that people (at least most people) even forgave me for not mentioning their own work.

It is possible to give a brief “coarse-grained” summary of the whole conference. This also is very easy, because it doesn’t vary from one conference to another. Here it is:

1. The Standard Model agrees with all data, but has many holes and loose ends.

2. We must still:
• find the top quark,
• identify the mechanism (or mechanisms) for CP nonconservation,
• solve quantum chromodynamics.

3. We don’t understand:
• why the parameters of the standard model take the values we observe,
• why there are three generations of quarks and leptons,
• why $SU(3) \otimes SU(2) \otimes U(1)$?

4. Most of these questions revolve around the central unresolved issue concerning the Standard Model: how is $SU(2) \otimes U(1)$ broken? There are two possibilities: The Goldstone bosons which provide the longitudinal parts of the W and the Z are either elementary, or they are not.

• If the Goldstone bosons are part of a multiplet of elementary scalars, then in order to explain why they are so light compared to a really fundamental scale we probably need to assume supersymmetry.

• If the Goldstone bosons are composites, then definitely there must be new extra-strong forces, like the technicolor forces reviewed here by Appelquist.
(Peskin in his talk referred to these alternatives as the weak-interaction route and the strong-interaction route.)

5. In order to resolve the issue of electroweak symmetry breaking and break out of the present impasse, we need the

- SSC
- LHC.

* * *

In the remaining 57 minutes of my talk I would like to discuss a few specific topics. These are chosen, not necessarily because they represent the most important things discussed at this conference, but mostly because they are matters that I found interesting and about which I had something that I wanted to say. In 1986 the two topics I chose to discuss were solar neutrino oscillations and string theory. (By the way, Robertson\footnote{I will follow the practice here for the most part of just quoting talks given at this conference, chiefly the rapporteur talks. References to the physics literature can be found in the written versions of these talks.} said in his talk the other day that only recently have particle physicists regarded the solar neutrino problem as part of particle physics, which is true only if “recently” extends back at least six years.) These would be a pretty good choice of a pair of special topics for my talk today (which shows how slowly our field is moving), and I will come back to them, but there are a few other
topics I want to take up here also. At the end of this discussion of special
topics I will say a little about what I think lies ahead for particle physics.

Heavy Quark Symmetry

My first topic is heavy quark symmetry. This was discussed by Drell and
Isigur in their rapporteur’s talks; summaries were given in parallel sessions
by Close and Grinstein; and the subject kept coming up in many of the talks
in parallel sessions I attended. It is really one of the prettiest developments
in the theory of strong interactions in a long time, and one in that in a
remarkable way has almost immediately been taken over by experimentalists
as part of their standard bag of tricks. Heavy quark symmetry has been
explained many times in this conference in a hand-waving way in terms of
analogies with atomic physics. Here I would like to offer an explanation
(most of which can be found in the papers of the experts), that you can give
with your hands strapped to your sides, in terms of the Feynman diagrams
of quantum chromodynamics.

Consider a heavy quark line going through a Feynman diagram. Suppose
that after emitting momentum to and absorbing momentum from the gluons,
it has acquired a four-momentum $mv + k$, where $m$ and $v$ are the mass and
initial four-velocity ($v \equiv [\gamma \vec{v}, \gamma]$) of the heavy quark. The components of the
four-momentum transfer are limited by an amount of the order of the QCD
scale factor $\Lambda$, so for $m \gg \Lambda$, the quark propagator may be approximated by

\[
\frac{-i(m \not{v} + k) + m}{(mv + k)^2 + m^2} \simeq \frac{-i \not{v} + 1}{2v \cdot k}.
\] (1)

[I am using the usual slash notation, $\not{a} \equiv \gamma^\mu a_\mu$, with a metric that has diagonal components $-1, +1, +1, +1$.] This immediately reveals the flavor degeneracy of these theories; as long as we express everything in terms of velocities rather than four-momenta, nothing in matrix elements depends on the heavy quark mass. With $N$ heavy types of heavy quarks, the bound states of one or more heavy quarks plus any number of light quarks and gluons would form $SU(N)$ multiplets of hadrons with equal binding energies.

To go further, suppose gluon lines with polarization indices $\mu$, $\nu$, etc. and color indices $a$, $b$, etc., are emitted by a heavy quark line before it finally leaves the diagram with spin $z$-component $\sigma$ and four-velocity $v$. By moving all factors $(-i \not{v} + 1)/2$ to the end of the line, we easily see that the contribution of the quark-gluon vertices and quark propagators is

\[
\bar{u}(\sigma, v) \left(\frac{-i \not{v} + 1}{2}\right) it_a \gamma^\mu \left(\frac{-i \not{v} + 1}{2}\right)
\times it_b \gamma^\nu \left(\frac{-i \not{v} + 1}{2}\right) \cdots
= \bar{u}(\sigma, v) \left(\frac{-i \not{v} + 1}{2}\right) t_a t_b \cdots v^\mu v^\nu \cdots .
\] (2)

In other words, nothing depends on the spin $z$-components of the heavy quarks, aside from kinematic final (and initial) state factors like $\bar{u}(\sigma, v)(-i \not{v} + 1)/2$, which do not vary from diagram to diagram for any
given process. (Despite appearances, this result does not depend on the fact that the quarks have spin \( \frac{1}{2} \)) In particular, the positions of the poles in a Feynman diagram will not depend on the spin \( z \)-components of any incoming or outgoing heavy quarks, so the bound states of heavy and light quarks will exhibit a spin degeneracy as well as a flavor degeneracy: all of the bound states that are related by rotating only heavy quark spins (as for instance the lowest 0\(^{-}\) and 1\(^{-}\) bound states of a heavy quark and a light antiquark) will have the same mass.

One can also apply this diagrammatic analysis to processes in which a weak interaction induces a transition between hadrons containing different species of heavy quarks. According to the foregoing analysis, the matrix element of a current \( \bar{\psi}_{f'} \Gamma \psi_f \) (with \( \Gamma \) an arbitrary \( 4 \times 4 \) matrix) between initial and final bound states \( \alpha \) and \( \alpha' \), consisting of any number of light quarks and/or antiquarks plus one heavy quark respectively of flavor \( f \) and \( f' \) and four-velocity \( v \) and \( v' \), must take the form:

\[
\langle f', v', \alpha' | \left( \bar{\psi}_{f'} \Gamma \psi_f \right) | f, v, \alpha \rangle = \sum_{\sigma', \sigma} C_{\alpha \rightarrow \alpha'}(\sigma', \sigma, v, \sigma) \times \left[ \bar{u}(\sigma', v') \left( -i \frac{g' + 1}{2} \right) \Gamma \left( -i \frac{\not{v} + 1}{2} \right) u(\sigma, v) \right],
\]

where \( C_{\alpha \rightarrow \alpha'} \) is an unknown but flavor-independent function of initial and final heavy quark spin \( z \)-components and velocities, that arises from the convolution of the wave functions for states \( \alpha \) and \( \alpha' \) with the sum of Feynman
diagrams in which the heavy quark line is replaced by a product of factors $v^\mu t_a$ for each interaction of gluons with the heavy quark. (These Feynman diagrams could actually be summed explicitly for Abelian gluons, but as far as is known not for real SU(3) gluons.) We can usefully rewrite this formula as a trace:

$$\langle f', v', \alpha' | \bar{\psi}_{f'} \Gamma \psi_f \rangle | f, v, \alpha \rangle = \text{Tr} \left\{ M_{\alpha \rightarrow \alpha'}(v', v) \left( -\frac{i \not{v} + 1}{2} \right) \Gamma \left( -\frac{i \not{v'} + 1}{2} \right) \right\},$$

where $M$ is the $4 \times 4$ matrix:

$$M_{\alpha \rightarrow \alpha'}(v', v) \equiv \sum_{\sigma, \sigma'} C_{\alpha \rightarrow \alpha'}(v', \sigma', v, \sigma) u(\sigma, v) \bar{u}(\sigma', v').$$

Though $M$ is unknown, its matrix structure is fixed by Lorentz invariance. For instance, for a transition between the lowest $0^-$ mesons with flavor $f$ and $f'$, the above matrix element must have the same Lorentz transformation properties as the matrix $\Gamma$, so $M_{0^- \rightarrow 0^-}(v', v)$ must be a scalar. Factors of $\not{v}$ and $\not{v}'$ do not matter because when multiplied into $(-i \not{v} + 1)$ or $(-i \not{v'} + 1)$ they merely yield factors of $+i$. Thus $M$ must here be proportional to the unit matrix:

$$M_{0^- \rightarrow 0^-}(v', v) = \frac{1}{2} \xi(-v \cdot v') 1,$$

with a coefficient $\xi$ that depends on the only scalar variable, $v \cdot v'$. This is the celebrated Isgur-Wise function. If you like analogies, then think of
\( \xi(-v \cdot v') \) as analogous to the well-known Fermi function \( F(Z,W) \), which gives the effect of final state Coulomb interactions in nuclear beta decay. The fact that \( \xi(-v \cdot v') \) depends neither on heavy quark flavor nor on the matrix \( \Gamma \) is just like the fact that \( F(Z,W) \) does not depend on the nature of the nuclei participating in the beta transition [aside from the energy \( W \), which is analogous to the variable \( v \cdot v' \), and the atomic number \( Z \), which is analogous to the fixed color triplet assignment of quarks], or on the matrices \( [S, V, T, A, \text{ or } P] \) appearing in the beta decay Hamiltonian.

It is a challenge to quantum chromodynamics to calculate the Isgur-Wise function, and much effort has been put into this problem, but there is one point where the value of \( \xi \) can be obtained without effort. In the special case where \( \Gamma = i\gamma^\mu \), \( f = f' \), and \( v = v' \), the matrix element (4) is given by the conservation of heavy quark flavor as just \( v^\mu \) [in much the same way that the matrix element for \( O^{14} \) beta decay is fixed by the conservation of the isospin vector current], so in this case \( \xi = 1 \) at the point \( v \cdot v' = -1 \). But \( \xi \) is a universal function, independent of \( \Gamma \) and heavy quark flavors, so in general

\[
\xi(1) = 1 .
\]

(7)

This result is used in determinations of CKM matrix elements like \( V_{cb} \) by extrapolating data on weak processes such as \( B \rightarrow D + \ell + \nu \) to the point \( v \cdot v' = -1 \).

The same Isgur-Wise function enters into the matrix elements for transitions involving the \( 1^- \) mesons that are degenerate with the lightest \( 0^- \)
mesons. Whatever the wave function for the lowest $0^{-}$ mesons, the heavy quark spin degeneracy tells us that the wave function for the degenerate $1^{-}$ meson [in its rest frame] with $J_z = 0$ is given by inserting an extra factor $(-1)^{2\sigma}$. We must therefore take $M_{0^{-}\to 1^{-}}$ as a scalar matrix function of the final vector meson polarization $e^\mu$ that in the rest frame of the final heavy quark for vector meson polarization vector $e^0 = e^1 = e^2 = 0, e^3 = 1$ gives a factor $\frac{1}{2}(-)^{2\sigma'}\xi$ when acting to the left on $\bar{u}(v',\sigma')$. Aside from inconsequential terms involving $\not{v}$ and/or $\not{v}'$, the unique matrix satisfying these requirements is

$$M_{0^{-}\to 1^{-},e}(v',v) = \frac{i}{2} \xi(-v \cdot v') \gamma_5 \gamma_\tau'.$$

(8)

This can be used in Eq. (4) to calculate the matrix elements for processes like $B \to D^\ast(2010) + \ell + \nu$. Similar formulas have been derived for $1^{-}\to 0^{-}$ and $1^{-}\to 1^{-}$ transitions, and for processes involving other hadrons containing one or more heavy quarks.

**High Precision Electroweak Physics**

Due largely to the great recent success of LEP, electroweak physics has become closer and closer in spirit to quantum electrodynamics, as a branch of physics where high precision is expected. As Rubbia said this morning, we now have three parameters in electroweak physics that we know with quite high precision: the Fermi coupling constant $G_F$ of beta decay (known from muon decay); the fine structure constant $\alpha$; and the mass of the $Z$,
measured at LEP. These three constants (along with masses and mixing angles for whatever quarks or leptons are involved in a given process) are all you need to calculate any desired matrix elements in the electroweak theory in tree approximation. Beyond the tree approximation one needs to know other parameters, like the top quark and the Higgs mass, but since these enter in loops observable quantities are less sensitive to them. As discussed by Rolandi, with our knowledge of $G_F$, $\alpha$, and $m_Z$ we are now able to use the electroweak theory to do high precision calculations of other measured quantities, such as the W mass, the Z leptonic width $\Gamma_\ell$, and $[\sin^2 \theta]_Z$ (measured from the forward-backward asymmetry in $e^+ + e^- \rightarrow \ell^+ + \ell^-$ at the Z peak), aside from a weak dependence on the top and the Higgs mass. These predictions agree with existing data for a top quark mass between 130 and 170 GeV (to 68 percent confidence) and a Higgs mass which is essentially unconstrained, allowed to be anywhere from 50 Gev to 1 TeV.

Satisfactory as this situation is, we can be pretty sure that it will “soon” (i.e., within a few years) be radically improved. We can count on a further improvement in the accuracy with which $m_W$, $\Gamma_\ell$, and $[\sin^2 \theta]_Z$ are known, especially through more accurate measurements of $m_W$ at LEP2 and the Tevatron Collider, and perhaps also through an improvement in the measurement of $[\sin^2 \theta]_Z$ at the SLD. But the real change will come when the top quark is discovered and its mass is measured, presumably within a few years at the Tevatron Collider. At that point we shall find ourselves in the position
of having a critical test of the simplest (one scalar doublet) version of the electroweak theory, and if that test is failed, of being able to say something about what new physics must be added to this model.

In the last few years it has become customary to parameterize the new physics that may enter in the electroweak theory in terms of what are called oblique radiative corrections. Parameters like the $Z$ mass, the $Z$ leptonic width, etc., would be affected by the top and Higgs mass, as well as by most kinds of new physics that could be added to the minimal electroweak theory, mostly through the $2 \times 2$ matrix vacuum polarization of the $Z$ and $\gamma$. (For instance, Higgs scalars do not interact very much with $u$ or $d$ quarks or electrons or muons because these quarks and leptons are so light.) Further, as Peskin explained, because experiments are still mostly done at energies that are rather small compared to what we think are the energy scales of new physics, the vacuum polarization can be parameterized by the coefficients of just the first few terms in a Taylor expansion. The most important of these coefficients are called $S$ and $T$. High precision measurements of $m_W$, $\Gamma_\ell$, and $[\sin^2 \theta]_Z$ define narrow strips of allowed values in the $S-T$ plane, all running with various slopes from $S < 0, T < 0$ to $S > 0, T > 0$. For a given top quark mass (measured, say, to an accuracy of 10 GeV) and a Higgs mass between 50 GeV and 1 TeV, the minimal standard model defines a short wedge in the $S-T$ plane, running roughly transverse to these strips. [These strips and wedges are shown in Figure 1. The strips in this figure are centered on
values of $S$ and $T$ derived from the current experimental values of $m_W$, $\Gamma_\ell$, and $[\sin^2 \theta]_Z$.

There are two possibilities for what will then be found.

- If the wedge intersects all the strips then, depending on where along the wedge the intersection is, we will have a good rough estimate of the Higgs mass in the minimal standard model. (It is often said that observables are only very weakly dependent on the Higgs mass in the minimal standard model, but that is true only if the top quark mass is not known; predictions for a low Higgs mass and a low top quark mass resemble those for a high Higgs mass and a high top quark mass.) If the experimentally allowed ranges of $m_W$, $\Gamma_\ell$, and $[\sin^2 \theta]_Z$ continue to be centered on their present values, then this intersection would indicate a relatively high Higgs mass for $m_t = 160 \pm 10$ GeV, and a relatively low Higgs mass for $m_t = 120 \pm 10$ GeV.

- On the other hand, if the wedge doesn’t intersect the strips, we will have a good clue as to what new physics must be added to the minimal electroweak theory. For instance, technicolor theories tend to move the wedges to larger values of $S$.

All this goes to underline the extreme importance of finding the top quark.
Effective Field Theories

My third topic is not new, but has lately become part of the common language of elementary particle physics. I think that I have heard the words “effective field theory” or “effective Lagrangian” a hundred times in as many different contexts at this meeting. Leutwyler went into effective field theories in some detail, and Peskin applied them in discussing oblique radiative corrections. I remember that after Leutwyler’s talk, there was a question from the audience asking how we know that it is legitimate to use these effective field theories in describing the real world, and then someone else came up to me in the lobby and asked the same question. So although this topic is not by any means new, I would like to take a few minutes to explain what we are doing when we use an effective field theory.

For this purpose I would like to use an example that is not usually discussed in terms of effective field theories, although I think it’s the first example of an effective field theory in the literature. It goes back to the 1930’s, when theorists like Heisenberg were calculating the scattering of light by light as a somewhat academic application of the new quantum electrodynamics. In 1936 H.Euler showed that the results for photon-photon scattering amplitudes at photon energies $\omega \ll m_e$, which are obtained in quantum electrodynamics from what would today be called an electron loop graph, could be summarized as a lowest-order perturbation theory result obtained from the Lagrangian density:
\[ L_{\text{eff}} = \frac{1}{2} (\vec{E}^2 - \vec{B}^2) + \frac{e^4}{360\pi^2m_e^4} \left[ (\vec{E}^2 - \vec{B}^2)^2 + 7 (\vec{E} \cdot \vec{B})^2 \right]. \] (9)

In modern terms, we would say that in order to study processes at energies much less than \( m_e \), we ‘integrate out’ the electron, replacing the Lagrangian of quantum electrodynamics with the effective Lagrangian (9).

This historical example provides lessons that help us to understand modern effective field theory:

1. We do not really need perturbation theory or even quantum electrodynamics to understand the general form of Eq. (9). Gauge and Lorentz invariance tell us that the most general possible Lagrangian for the electromagnetic field is of the form

\[ L_{\text{eff}} = \frac{1}{2} (\vec{E}^2 - \vec{B}^2) + c_1 (\vec{E}^2 - \vec{B}^2)^2 + c_2 (\vec{E} \cdot \vec{B})^2 + c_3 (\vec{E} \cdot \vec{B}) \Box (\vec{E} \cdot \vec{B}) + \cdots, \] (10)

where the ellipsis “\( \cdots \)” denotes other terms with derivatives and/or more electromagnetic field factors, and \( c_1, c_2, c_3, \text{etc.} \), are unknown constants. [The factor in the first term is made to be equal to \( \frac{1}{2} \) by a canonical normalization of the fields.] Dimensional analysis tells us that \( c_1 \) and \( c_2 \) have dimensionality \( \text{[mass]}^{-4} \), while \( c_3 \) and all other coefficients have dimensionalities \( \text{[mass]}^{-n} \), with \( n \geq 6 \). If (10) is obtained by integrating out ‘heavy’ charged particles (like the electron)
with some typical mass $M$, then $c_1$ and $c_2$ will be proportional to $M^{-4}$, while $c_3$ and all other coefficients will be proportional to $M^{-6}$ or higher powers of $1/M$. If we calculate photon scattering amplitudes at photon energies $\omega \ll M$, then the result will be dominated by the Born approximation contribution of just the $c_1$ and $c_2$ terms in (10), because all other interactions and higher order graphs will be suppressed by two or more powers of $\omega/M$. [Of course one must go back to quantum electrodynamics to derive the specific values of the coefficients in Eq. (9), though factors like $e^4$ and $1/8\pi^2$ can be obtained by simply counting vertices and loops in the graphs from which the effective Lagrangian is calculated.]

2. The effective Lagrangian (10) is more than an aide memoire for photon-photon scattering amplitudes. It is not hard to show that the dominant terms in the amplitudes for other low-energy photon reactions like $\gamma + \gamma \rightarrow \gamma + \gamma + \gamma + \gamma$ are given by using the $c_1$ and $c_2$ terms in (10) in the tree approximation. Perhaps more surprisingly, we can regard this effective Lagrangian as the basis of a legitimate quantum field theory, calculating corrections to any process of higher order in $\omega/M$ by including photon loop as well as tree diagrams generated by (10). This effective field theory is of course non-renormalizable, but the divergences in these loop graphs are cancelled by renormalization of the coefficients in (10). Non-renormalizable theories are just as renormalizable as renormalizable
theories; the only difference is that we have to deal with an infinite number of interaction types. You might suppose that the intrusion of unknown constants from the higher terms in (10) means that this effective theory has lost all predictive power, but that’s not true at all. For example, the one-loop contribution to the scattering of light by light gives a term of order $(\omega/M)^8 \log \omega$ with a coefficient given by a known quadratic expression in $c_1$ and $c_2$, as well as polynomial terms involving the coefficients of higher terms in (10).

3. The use of perturbation theory with the Lagrangian (10) does not depend on the fact that the fine structure constant is small (though that helps); perturbation theory could be used even if the underlying theory, quantum electrodynamics, were a strongly interacting theory. This is because each loop in the effective field theory introduces additional factors of $\omega/M$ into the amplitude for any given process; perturbation theory here is an expansion in powers of $\omega/M$, as well as $e^2/8\pi^2$. This of course is why Euler was able to calculate the leading terms in the photon-photon scattering amplitude by using (9) in lowest order. If $e^2/8\pi^2$ were of order unity we would not be able to derive Eq. (10) from an underlying theory, and we would not be able to calculate values for the constants $c_1$, $c_2$, etc., but we could still do perturbation theory with (10) as our effective Lagrangian.

4. The crucial feature of the effective Lagrangian (10) that allows us to
use it to generate an expansion for amplitudes in powers of energy is that it involves only non-renormalizable interactions, with coupling constants that necessarily have the dimensions of negative powers of mass. In pure electrodynamics this is a consequence of gauge invariance; renormalizable interactions would be quartic polynomials in the vector potential $A^\mu$ without derivatives, and these would not be gauge invariant.

The last remark suggests that we should expect to be able to use effective Lagrangians to do calculations at low energies whenever there are symmetries (like gauge invariance in the above example) that rule out renormalizable interactions. There are other such symmetries. One is general covariance in the theory of gravitation (with no cosmological constant.) Another well-known example is the broken chiral $SU(2) \otimes SU(2)$ symmetry of quantum chromodynamics with massless $u$ and $d$ quarks. This requires that the pion, the Goldstone boson of this broken symmetry, enters into the Lagrangian only in non-renormalizable couplings involving derivatives of pion fields, such as the leading term in the purely pionic Lagrangian:

$$L_{\text{eff}} = - \frac{\partial_\mu \vec{\pi} \cdot \partial^\mu \vec{\pi}}{(1 + \vec{\pi}^2/F_\pi^2)^2}.$$  \hspace{1cm} (11)

Effective Lagrangians were introduced into modern particle physics in the 1960s in this context (though we did not know we were doing quantum chromodynamics then.) One recent development in this area is the extension of
the effective Lagrangian method to not only the interactions of pions with each other and with single nucleons, but also to the interactions of pions and several nucleons - in other words to the problems of nuclear force and pion scattering on nuclei. The results obtained can be summarized by saying that the chiral Lagrangian approach turns out to justify approximations (such as assuming the dominance of two-body interactions) that have been used for many years by nuclear physicists, though not knowing about the chiral Lagrangian approach it was not quite fair of them to use these approximations.

The heavy quark symmetries discussed earlier have been incorporated into an Effective Heavy Quark Field Theory, and this has been combined with the soft pion theory based on broken chiral symmetries. Effective Lagrangian techniques are also being used today to study not only the ordinary strong interactions, but also technicolor, a hypothetical extra strong force that has been considered as a possible source of the breaking of the $SU(2) \otimes U(1)$ electroweak symmetry. It is easy to invent technicolor models that have an accidental $SU(2) \otimes SU(2)$ symmetry containing the $SU(2) \otimes U(1)$ symmetry of the electroweak interactions, and it’s natural to assume that it is spontaneously broken to a sort of isospin symmetry called custodial $SU(2)$, which preserves the usual relation between $W$ and $Z$ masses. This symmetry breakdown of course entails Goldstone bosons that provide the longitudinal parts of the $W$ and the $Z$, and in most cases additional Goldstone bosons called techni-pions. The effective Lagrangian of these Goldstone bosons is very
much like Eq. (11), with a structure dictated by the broken and unbroken symmetries and by the nature of the degrees of freedom. As discussed here by Chanowitz, effective Lagrangians of this sort are actively being used to study the possibilities of finding signs of technicolor at accelerators like the SSC.

This brings us back to the question that was asked after Leutwyler’s talk: how do we know that it is legitimate to use these effective Lagrangians? The derivation of a Lagrangian like (11) is not just a matter of ‘integrating out’ heavy degrees of freedom in quantum chromodynamics, in the way that Euler integrated out the electron to obtain (9), because unlike the electron in QED, the pion does not appear as a field in the QCD Lagrangian. True, the breaking of the \( SU(2) \otimes SU(2) \) symmetry implies that there must be a pion particle, and dictates the structure of any field theory that describes this pion, but how do we know that this composite pion can be described by a field theory at all? The answer seems to be that any quantum theory that satisfies Lorentz invariance plus a technical requirement called cluster decomposition plus unitarity will always at sufficiently low energy look like a quantum field theory. Quantum field theory is the only way, we believe, of reconciling these fundamental requirements. [The clause “at sufficiently low energy” is inserted so that this statement will apply also to string theories, which contain infinite numbers of increasingly heavy particle types, and therefore do not look like quantum field theories at energies comparable to the string

19
mass scale.] Furthermore, the theory of low energy pions must be a quantum field theory described by a broken $SU(2) \otimes SU(2)$ symmetry. The use of Eq. (11) is justified not because we derive it from a field theory like QCD, but because it is the most general chiral invariant theory of pions, aside from terms with more derivatives whose effects are suppressed at energies $E$ much less than a QCD scale of order $m_\rho$ by two or more factors of $E/m_\rho$.

The same thought occurs to us with somewhat chilling overtones with regard to the Standard Model. We believe that there is some unknown physics at really high energies, roughly of the order of the Planck mass, with unbroken symmetries like $SU(3) \otimes SU(2) \otimes U(1)$ (and perhaps supersymmetry) that keep the particles of the standard model massless before the breaking of $SU(2) \otimes U(1)$ (and supersymmetry). The most general theory we could expect to find at low energies is simply the most general quantum field theory satisfying these symmetries. This is of course the standard model, supplemented with non-renormalizable terms whose effects at energy $E$ are suppressed by powers of $E/m_{\text{Planck}}$. The success of the standard model tells us nothing about whether the underlying theory that describes physics at the Planck scale is a quantum field theory. In a sense this represents the revenge of S-matrix theory, because we now believe that the field theories of which we are so proud, quantum electrodynamics, quantum chromodynamics, even general relativity for that matter, are not what they are because they are truly fundamental field theories, but simply because they are the only way
of satisfying the requirements of symmetries and S-matrix theory. Quantum
field theory as we use it now is nothing but S-matrix theory made practical.
Our best candidate for an underlying theory at the Planck scale is in fact not
a quantum field theory in the ordinary sense, which brings me to my next
topic.

**Superstring Theory**

Superstring theories have been studied for over a decade as candidates
for a fundamental theory in particle physics. The implications of superstring
theories were discussed in the rapporteur talks by Alvarez Gaumé and Pe-
skin, and reviewed here in detail in a parallel session talk by Dine; in this
summary I will just make some general remarks about the current state of
the theory, with emphasis on recent work on coupling constant unification in
these theories.

One change in the last few years has been a movement away from thinking
of superstrings as existing in ten or twenty-six dimensions, of which all but the
four dimensions of ordinary spacetime somehow become compactified, and
toward formulating superstring theories from the beginning in four spacetime
dimensions. Not that the previous theories are wrong — it’s just that if one
starts with a ten dimensional superstring theory and then supposes that six
of these dimensions become compactified, what you get is the same as if you
had started with a four dimensional superstring theory of an appropriate
type. So the four-dimensional approach gives a more general approach to theories that have a chance to describe nature.

A string as it moves through space sweeps out a two dimensional surface, so the study of strings is the study of quantum field theories in two dimensions, with the spacetime coordinates along with other degrees of freedom appearing as fields in these theories. The two surface coordinates in these field theories may be represented as a single complex variable $z$, but with the complex plane given an arbitrary topology, one that becomes more and more complicated as we go to higher and higher order in perturbation theory.

The action for this two-dimensional theory is a functional of the four spacetime coordinates $x^\mu(z, z^*)$, plus [in the popular ‘heterotic’ superstring theories] an equal number of spinor coordinates $\psi^\mu(z)$, plus a number of other field variables that are needed to satisfy certain constraints. These constraints arise ultimately from the fact that the spacetime metric has $g_{00} = -1$. This would destroy the positive definiteness of the quantum field theory of the coordinates $x^\mu(z, z^*)$ if it were not for a symmetry known as conformal invariance, which allows us to transform away vibrations of the string into the timelike direction. But conformal invariance in a theory with just four $x^\mu(z, z^*)$’s or four $x^\mu(z, z^*)$’s and four $\psi^\mu(z)$’s is violated by quantum mechanical anomalies similar to the triangle anomaly in QCD, and these anomalies must be cancelled by the other fields added to the action.

Each possible equilibrium state of ordinary fields in four dimensions corre-
sponds to a different two-dimensional conformal field theory, and the vacuum expectation values of the four-dimensional fields correspond to parameters in the conformally invariant action. For instance, there is a term in the string action that (after a suitable choice of coordinates in the complex plane) is of the form:

\[
I_{quad}[x] = -\frac{1}{2} \int \! \! \int dz \, dz^* \, g_{\mu\nu}(x(z, z^*)) \times \frac{\partial x^\mu(z, z^*)}{\partial z} \frac{\partial x^\nu(z, z^*)}{\partial z^*},
\]

where \( g_{\mu\nu}(x) \) is the gravitational field, which is required by conformal invariance to satisfy the Einstein field equations.

There are some useful general results that apply to this whole class of theories in the tree approximation, where the complex plane is not given any topological complications. One result is a formula for the string mass, a parameter related to the string tension and also to the slope \( \alpha' \) of the Regge trajectories on which the excited states of the string lie, and which can most simply be regarded as the mass of the first excited state of a vibrating string. It is given in terms of the string coupling constant (which determines the magnitude of higher order effects) and the Planck mass by

\[
M_{\text{string}} \equiv \frac{2}{\sqrt{\alpha'}} = g_{\text{string}} \frac{M_{\text{Planck}}}{\sqrt{8\pi}}.
\]

Also, if physics at energy scales below \( M_{\text{string}} \) is described by an \( SU(3) \otimes SU(2) \otimes U(1) \) gauge field theory, then the gauge coupling constants at a
renormalization scale $\mu$ near $M_{\text{string}}$ approach values satisfying the relation:

$$k_{SU(3)} g_{SU(3)}^2 = k_{SU(2)} g_{SU(2)}^2 = k_{U(1)} g_{U(1)}^2 = g_{\text{string}}^2.$$  \hfill (14)

The constants $k$ are known in the trade as the ‘levels’ of the Kac-Moody algebra from which the gauge symmetries are derived, but the important thing is that $k_{SU(3)}$ and $k_{SU(2)}$ are in general integers, in most conformal field theories just $+1$. Also, $k_{U(1)}$ takes discrete values (values that are not changed by continuous changes in the conformal field theory), and in a large class of interesting models takes the familiar value of $5/3$. In other words the relations among gauge couplings, that until recently have been understood in terms of the ‘grand unification’ of $SU(3) \otimes SU(2) \otimes U(1)$ in some simple group, occur naturally in string theory without any need for grand unification. In fact, string theory gives a better explanation of these coupling constant relations than grand unification because in grand unified theories the presence of color triplet partners of the Higgs doublet makes it particularly difficult to understand the disparity between the electroweak scale and the grand unification scale, while such GUT partners are absent in string theories with $k_{SU(3)} = k_{SU(2)} = 1$.

Eq. (14) applies to gauge couplings defined at a renormalization scale that is roughly of the order of the string mass scale (13). For more precise information about the value of the scale (or scales) where these relations apply, one must look to a one loop calculation. Such calculations were described here briefly by Alvarez Gaumé, and would have been described by
Peskin if time had allowed. Including one-loop corrections, the relation (14) is replaced by

\[ \frac{16\pi^2}{g_a^2(\mu)} = k_a \frac{16\pi^2}{g_{\text{string}}^2} + b_a \ln \left( \frac{M_{\text{string}}^2}{\mu^2} \right) + \Delta_a, \tag{15} \]

where \( b_a \) are the familiar constants appearing in the one-loop renormalization group equations \( \left[ \mu d g_a/d\mu = b_a g_a^3/16\pi^2 \right] \) and \( \Delta_a \) are the so-called threshold terms. In string theory as in field theory, the threshold terms are given by a sum over all species \( n \) of heavy particles that are integrated out to obtain the effective “low-energy” gauge theory:

\[ \Delta_a = \sum_n c_{na} \ln \left( \frac{M_n^2}{M_{\text{string}}^2} \right), \tag{16} \]

with \( M_n \) equal (up to factors of order unity) to the heavy particle masses, and \( c_{na} \) constants that depend on the details of the underlying theory. The \( c_{na} \) are of order unity, and we expect the \( M_n \) to be of the order of \( M_{\text{string}} \), so the threshold corrections \( \Delta_n \) may be expected to be of order unity, which is why we expect Eq. (14) to apply at a renormalization scale \( \mu \) of order \( M_{\text{string}} \).

To make a more precise estimate, it is convenient to rewrite Eq. (15) as:

\[ \frac{16\pi^2}{g_a^2(\mu)} = k_a \frac{16\pi^2}{g_{\text{string}}^2} + b_a \ln \left( \frac{M_{\text{string}}^2}{\mu^2} \right) + \tilde{\Delta}_a, \tag{17} \]

where

\[ \tilde{\Delta}_a \equiv \sum_n c_{na} \ln \left( \frac{M_n^2}{M_{\text{string}}^2} \right) + b_a \ln \left( \frac{M_{\text{string}}^2}{M^2} \right), \tag{18} \]
and $M$ is an arbitrary mass. The point is to try to choose $M$ to minimize the typical values of $\tilde{\Delta}_a$, in which case according to (17) $M$ is the renormalization scale at which the couplings most closely satisfy (14). This has been done by a number of authors; what seems to be the best estimate is:

$$M = \frac{e^{(1-\gamma)/2}3^{-3/4}}{\sqrt{2\pi}} \ M_{\text{string}} = 0.216 \ M_{\text{string}}.$$  

(19)

Extrapolating experimental values of the couplings to high energy in the minimal supersymmetric standard model, the couplings $g_{SU(3)}^2/4\pi$, $g_{SU(2)}^2/4\pi$, and $\frac{5}{3}g_{U(1)}^2/4\pi$ are found to converge to a common value $1/(26.3 \pm 2.1)$ at a renormalization scale about equal to $2 \times 10^{16}$ GeV. [It has recently become possible to be much more precise about this, largely because experimental and theoretical advances have produced a great improvement in our knowledge of $g_{SU(3)}^2$, reported in a plenary session here by Bethke.] With $g_{\text{string}}^2/4\pi$ taken equal to $1/26.3$, Eq. (13) gives $M_{\text{string}} = 1.7 \times 10^{18}$ GeV, and Eq. (19) then gives the renormalization scale for coupling constant unification as $M = 3.6 \times 10^{17}$ GeV. So in other words there is about a factor of twenty discrepancy between the unification scale expected in string theory, proportional to the Planck mass, and the value $2 \times 10^{16}$ GeV inferred from ‘low’ energy experiments. In speaking of this as a discrepancy, I can’t help feeling a sense of unreality. Here we are, talking about energy scales about $10^{13}$ to $10^{14}$ times larger than the largest energies that we can produce in our accelerators, and worried about a discrepancy of a factor of 20!

Of course, one possible resolution of this discrepancy is that there may
be more to physics at energies below the string scale than the minimal supersymmetric standard model, and that the couplings have therefore not been extrapolated correctly to very high energies. It is not easy to see what could be added to the minimal supersymmetric standard model that would preserve the natural convergence of all three couplings to a common value, while moving the energy at which the convergence occurs to higher values. Alternatively, the resolution of this discrepancy may be that the threshold corrections are larger than we had thought. From (16), we see that this could happen if some of the particles that were ‘integrated out’ in deriving the $SU(3) \otimes SU(2) \otimes U(1)$ standard model are actually somewhat lighter than $M_{\text{string}}$. Such would be the case if physics immediately below the string scale were described by a Kaluza-Klein theory, with some extra spatial dimensions forming a compact manifold with dimensions somewhat larger than $1/M_{\text{string}}$, or by a four-dimensional grand unified theory, with a simple gauge group broken spontaneously to $SU(3) \otimes SU(2) \otimes U(1)$ at an energy somewhat smaller than the string scale. But it now seems that if Kaluza-Klein or grand unified theories are to have any application to the real world, it will only be in a narrow energy range, roughly from $4 \times 10^{17}$ GeV down to $2 \times 10^{16}$ GeV.

These threshold corrections are important in other ways. It has been known for some years that parameters of the conformal field theory such as the dilaton field and modular fields can be fixed only by a non-perturbative dependence of the vacuum energy on these parameters, because in the range
of these parameters where perturbation theory is valid, one can show that
the vacuum energy has no local minimum. [The dilaton field is particularly
important, because it is directly related to the gauge couplings.] We know
that the vacuum energy depends non-perturbatively on the couplings $g_a^2(M)$;
for instance the coupling $g_a^2(\mu)$ of QCD or some ‘hidden sector’ gauge field
becomes strong at a renormalization scale $\mu_a$ of order
$$
\mu_a \approx M \exp \left( -\frac{8\pi^2 k_a}{g_a^2(M)} \right),
$$
and the vacuum energy density contains terms of order $\mu_a^4$. But (15) shows
that the couplings $g_a^2(M)$ depend on the threshold correction terms $\Delta_a$, which
depend in a calculable way on parameters like the dilaton and modular pa-
rameter fields, giving the vacuum energy the desired non-perturbative de-
pendence on these parameters.

This analysis has been explored by many authors in the last few years,
and there are now many candidates for conformal field theories in which by
‘discrete fine tuning’ (that is, by choosing specific models out of a list of
thousands, but without carefully adjusting continuous parameters) you can
make the string coupling come out to have the ‘observed’ value, $g_{\text{string}}^2/4\pi =
1/26.3$, and you can also arrange to have supersymmetry broken at a scale
that would give the gravitino a mass of order 1 TeV, needed to produce
electroweak symmetry breaking with the observed strength. This seems to
me to represent real progress in the interaction between superstring theory
and physics.

28
Speaking of progress, there has been progress of another sort lately. It has actually become easier to follow the superstring literature, because many superstring theorists are now working on quantum gravity or two-dimensional statistical mechanics, and therefore there is less to read that is relevant to particle physics.

Despite all this progress, superstring theory faces a number of important and difficult problems. One of them is to identify correctly the source of non-perturbative effects. The non-perturbative effects described above are found by studying the effective ‘low’ energy quantum field theory derived from string theory rather from string theory itself, but we do not know if these the only important non-perturbative effects in string theories. For some years it has been widely assumed that the way to get at other, really stringy, non-perturbative effects is through what’s called string field theory. String field theory allows for the creation and annihilation of strings in much the same way that ordinary field theory describes the creation and annihilation of particles. I have never been enthusiastic about string field theory, because it seemed to me to take a beautiful new formalism and make it ugly by trying to make it look like the old formalism of quantum field theory, but who knows?

I would like to offer a modest suggestion as to where we might look for specifically stringy non-perturbative effects without developing a string field theory. Let’s recall what Feynman diagrams signal the advent of non-
perturbative effects in ordinary quantum field theory. Suppose you want to calculate some process like a vacuum polarization in a pure non-Abelian gauge field theory, at a momentum $p$ that is very small compared to the string scale $M_{\text{string}}$. Instead of introducing running coupling constants, suppose we perversely continue to expand in the coupling defined at $M_{\text{string}}$ [essentially $g_{\text{string}}$], as if Gell-Mann and Low had never been born. We would then find a breakdown of perturbation theory, because small factors of $g_a(M_{\text{string}})^2/8\pi^2$ would be accompanied with large factors of $\ln(p/M_{\text{string}})$. These large logarithms come from graphs in which internal massless gauge boson lines approach the mass shell. Now, in string theory the Riemann surfaces that correspond to these diagrams have handles that are pulled out to long thin tubes. But just as in field theory the large logarithms mean that we have to deal with loops within loops within etc., in string theory at low momentum we have to deal with thin handles attached to thin handles attached to etc. The Riemann surfaces from which non-perturbative effects arise, though infinitely complicated, may like a fractal surface have a self-similarity property, of looking qualitatively the same at all scales.

I have no idea how to deal with fractal Riemann surfaces, but it may be simpler than the task of dealing with the corresponding non-perturbative effects in quantum field theory, because as Alvarez Gaumé emphasized, in some respects string theory is much simpler than quantum field theory. Gauge invariance appears more naturally, and there are fewer diagrams. This aspect of
string theory has been exploited as a calculational device in quantum chromodynamics, and may perhaps allow us to understand stringy non-perturbative
effects that go beyond anything we have been able to understand in quantum
field theory.

Perhaps the most fundamental problem facing string theory is that we
still do not know even in principle what it is that chooses the correct string
theory, corresponding to the correct vacuum. The wrong answer is that the
correct vacuum is the one with lowest vacuum energy, because we already
know plenty of superstring theories that have negative vacuum energy. There
is another possible answer supplied by quantum cosmology. In recent years
the study of wormhole effects has suggested that the universe is not in a state
in which all coupling constants have definite values, but rather in a quantum
mechanical superposition of such states. Some relations among coupling con-
stants are fixed by fundamental principles (such as the fact that the electric
charges of the electron and positron are equal and opposite) and would be the
same in all terms in this superposition, but any coupling constant that can
vary continuously from one theory to another would have a continuous range
of values in the different terms in the state vector of the universe. There may
not be any such free parameters, in which case the following discussion is ir-
relevant, but our failure to formulate any principle for choosing the vacuum
state in string theories suggests that the vacuum energy may be a free pa-
rameter that varies continuously from term to term in the state vector of the
universe. When the universe becomes large, as it is now, it begins to look like
an incoherent mixture of these terms, with various probabilities. In this case,
it is only common sense that scientists who worry about the vacuum energy
would have to find a value in the fairly narrow range in which life could arise;
all the other terms in the wave function are there, but there are no scien-
tists to observe them. The vacuum energy acts in cosmology like Einstein’s
cosmological constant. It can’t be too positive because then galaxies would
never have formed, and it can’t be too negative because then the universe
wouldn’t live long enough for life to evolve. On this basis we can understand
in a natural way why the vacuum energy is relatively small - some hundred
and twenty orders of magnitude smaller than we would guess (a Planck mass
per cubic Planck length) from purely dimensional considerations. But these
considerations do not tell us that the vacuum energy (or equivalently, the
cosmological constant) is zero, or even that it is astronomically negligible,
but only that it is less than about a hundred times the present mass density
of the universe. This is an interesting bound, because both the cosmological
missing mass problem and the cosmic age problem mentioned by Krauss (the
fact that age of globular clusters seems to be larger than some estimates of
the age of the universe) would be alleviated by a positive cosmological con-
stant corresponding to a vacuum energy roughly ten times the present mass
density of the universe. It will be interesting to see whether this is the case.

* * *

32
This concludes my survey of special topics. I will now stick my neck way out, and try to guess what lies ahead for particle physics.

• All the present experimental challenges to the standard model will disappear. This judgment is based in part on having lived through it all before. Experiments certainly from time to time have required the incorporation of new features in the standard model, such as the tau lepton and the bottom quark. But where new experimental results have proved indigestible, irreconcilable with the general framework of the standard model, they have always gone away. Among these supposed difficulties were the high-$y$ anomaly, parity conserving neutral currents, anomalous trimuon events, and second-class currents. (Some of us suffer just hearing this list.) The two outstanding present experimental difficulties with the standard model are the tau branching ratios discussed by Drell and the seventeen kilovolt neutrino discussed by Robertson. Drell’s talk hints that the tau branching ratio problem is beginning to go away (although it certainly hasn’t gone away yet), and Robertson came out pretty strongly against the reality of the seventeen kilovolt neutrino. So it takes no great courage to predict that these anomalies will go away as well.

• My second guess (these are all just guesses) is that the electroweak symmetry will turn out to be broken by the vacuum expectation val-
ues of elementary scalars that appear in the effective Lagrangian at accessible energies, like the scalar doublet in the original electroweak theory, and that the hierarchy problem will be solved by supersymmetry. I say this for a number of reasons. As Peskin discussed, the technicolor idea requires awkward extensions to produce the quark and lepton masses without introducing new difficulties like flavor changing neutral currents. Also, as Peskin and Rubbia remarked, some simple technicolor theories are already excluded by the high precision electroweak data discussed above. Another reason for this guess is that I find the convergence of the $SU(3) \otimes SU(2) \otimes U(1)$ couplings in the supersymmetric standard model very impressive, and this convergence is easily lost if you mess up the model by adding things like technicolor. My last reason has to do with the solar neutrino problem, and is explained below.

- I would guess that the solar neutrino deficit is real and is in fact explained by the MSW effect. This is in part because of the present state of the neutrino experiments. As discussed by Krauss and in a parallel session by Bahcall, it is not possible by adjusting the temperature at the center of the sun to make the standard solar model fit both the chlorine and Kamiokande data, but calculations based on the MSW effect can fit everything, including the data from SAGE and GALLEX. Furthermore, from a theorist’s point of view the neutrino mass-square
difference $\Delta m_{\nu}^2$ and mixing angle $\theta_\nu$ that are needed to fit this data are very plausible: $\theta_\nu$ is like a typical small mixing angle in the CKM matrix, and $\sqrt{\Delta m_{\nu}^2}$ is a few millivolts, which is just what would be expected in the simplest extensions of the standard model. As is always the case for effective field theories, the Lagrangian of the standard model must be supplemented with non-renormalizable terms that are suppressed by powers of some large mass $M$, such as $10^{16}$ GeV. The least suppressed non-renormalizable term is a quartic term of dimension five, involving two factors of both the Higgs doublet and the lepton doublets:

$$L_5 = \frac{g_{ij}^2}{M} \left[ \left( \frac{\phi^0}{\phi^+} \right) \cdot \left( \frac{\nu_i}{\ell_i} \right) \right]$$

$$\times \left[ \left( \frac{\phi^0}{\phi^+} \right) \cdot \left( \frac{\nu_j}{\ell_j} \right) \right],$$

that after electroweak symmetry breaking yields a Majorana neutrino mass matrix:

$$m_{\nu ij}^2 = g_{ij} \langle \phi^0 \rangle_{\text{vac}}^2 / M.$$  \hspace{1cm} (21)

If we take the couplings $g_{ij}$ to be of the order of the product of the Yukawa couplings of the $i$’th and $j$’th lepton doublets to the scalar doublets, then the largest neutrino mass is of the order of $m_{\text{top}}^2 / M$, which for $M \approx 10^{16}$ GeV is indeed a few millivolts. But this attractive picture of the origin of the neutrino masses needed in the MSW effect is only possible if the scalars are elementary. If in a technicolor theory
you try to construct lepton number violating non-renormalizable interactions, you must replace the scalar doublet with some sort of bilinear function of techniquark fields, but then (21) is replaced with an operator of very high dimension, which is strongly suppressed by many factors of $1/M$. Of course, a neutrino mass of a few millivolts might arise from all sorts of possible new physics, like lepton symmetry breaking at the technicolor scale, but there would be no particular reason to expect millivolt neutrino masses. This gives a special importance to studies of solar neutrinos and neutrino oscillations, but the pace of these experiments is unfortunately very slow, a little like real time studies of continental drift. In the next few years we can look forward to SAGE and GALLEX being calibrated with artificial megacurie neutrino sources, to super Kamiokande coming into being, and to the start of the SNO experiment in Canada. All of these are important, but I want to emphasize that the SNO experiment offers the possibility of measuring a process $\nu + d \rightarrow \nu + p + n$ that arises solely from weak neutral currents and should therefore be unaffected by the MSW effect. If the MSW effect is indeed solely responsible for the observed neutrino deficit, then this neutral current process should be observed at precisely the rate predicted by the standard solar model, which may be our best way of reassuring ourselves that the sun is really well understood.

- We are going to find a great deal of new physics at accessible accelerator
energies. With supersymmetry invoked to solve the hierarchy problem, we expect not only sparticles, but also flavor changing neutral current processes that as Peskin emphasized are endemic in supersymmetry theories. Also endemic in supersymmetry theories are CP violations that go beyond the CKM matrix, and for this reason it may be that the next exciting thing to come along will be the discovery of a neutron or atomic or electron electric dipole moment. These electric dipole moments were just briefly mentioned at this conference, but they seem to me to offer one of the most exciting possibilities for progress in particle physics. Experiments here as in solar neutrino physics move very slowly, but I should mention that there has been a lot of progress lately in calculating the electric dipole moment of atoms in various models, with results that are encouraging for future experiments.

- The correct theory underlying the standard model is probably a superstring theory. So far, our best proof consists of asking what else it could be. Superstring theories may be confirmed (and here I'm saying something that Peskin especially wanted me to say) by predictions for the coefficients of soft supersymmetry breaking terms in the supersymmetric standard model. In particular, it has been recently realized that in superstring theories it’s typical that the lightest superparticles are gauginos rather than squarks or sleptons. This brings me to my final rash remark:
• Photinos are the cold dark matter needed in galaxy clusters.

* * *

In sticking my neck way out and acting as if we really are beginning to see the final answer I am going against the conventions of conference summarycraft. Michael Peskin showed a cartoon that expresses a more common view: a complacent physicist working away at a computer terminal, oblivious to monsters lurking just behind a wall. It is usual to end a conference summary with a remark that of course we expect that we will discover entirely new physics and that we are very far from anything like a final answer. But just because this is the conventional view and has generally been true in the past does not mean that it is true now. Michael mentioned a geographical metaphor: Columbus expecting to sail straight to the Indies, and not realizing there was something equally interesting in the way, namely America. [By the way, that geographical metaphor was used at a breakfast meeting in 1987 in convincing the Secretary of Energy to support the SSC.] But maybe a different geographical metaphor is more to the point. Imagine polar explorers sitting in the Travellers Club in the late nineteenth century, saying over their port, "You know, no matter how far north one goes, there’s always plenty of sea and ice left further north. No matter how far north we go, we shall never get to the North Pole." Well, eventually explorers did get to the North Pole. We too may eventually get to our destination, to a final theory, and possibly sooner rather than later.
Acknowledgement

I wish to thank John Bahcall, Michael Dine, Lance Dixon, Willy Fischler, Howard Georgi, Nathan Isgur, Vadim Kaplunovsky, Michael Peskin, and Joe Polchinski for helpful conversations regarding topics covered in this report.