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PHENOMENOLOGY BEYOND THE STANDARD MODEL

G.G. Ross^{*)}

CERN - Geneva

A B S T R A C T

A review is presented of recent theoretical and phenomenological studies of possible structure going beyond the standard model.

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*) On sabbatical leave from Department
of Theoretical Physics, University of
Oxford, 1 Keble Road, Oxford, U.K.

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1. INTRODUCTION

In view of the fact that this conference has largely been a celebration of the success of the standard model predictions, any review of progress in our attempts to foresee the physics that lies beyond the standard model should perhaps start with the reasons why it is widely believed that the standard model is not the final theory.

Of course, there are many aesthetic reasons to think that the unification of the fundamental forces including gravity should be achieved in a final "theory of everything". Certainly one hopes that such a theory should determine the numerous masses, couplings and multiplet structure of the standard model. However, even without these prejudices for Grand Unification, it is quite remarkable how closely the standard model resembles what we would expect of an effective field theory which would result if there is a scale Λ , beyond which new physics occurs. To see this note that in such a theory we would expect all our particle states to have a mass of order Λ unless there is a symmetry forbidding such a mass term from occurring. The light states of the standard model obey this rule (almost!). The spin-1 states of the theory are all associated with gauge bosons of the $SU(3) \times SU(2) \times U(1)$ gauge group and the gauge principle prevents them from acquiring a mass until this symmetry is spontaneously broken. Thus their mass is not of order Λ but of order the electroweak breaking scale. The spin- $\frac{1}{2}$ states of the theory, the quarks and the leptons, are in chiral representations, i.e., their left- and right-handed components transform differently under the gauge group. For this reason they, too, cannot acquire a mass for a (Dirac) mass term couples left- and right-handed components together and this cannot lead to a gauge invariant form. Once again the mass of these states is not governed by the scale Λ but by the electroweak breaking scale.

*) On sabbatical leave from : Department of Theoretical Physics, University of Oxford, 1 Keble Road, Oxford, U.K.

Thus the gauge structure of the vector bosons and chiral structure of the fermions in the standard model is just what we would expect from an effective field theory in which there is a new scale of physics, Λ , much higher than the electroweak breaking scale. The effects of this new physics will only be manifest through higher dimension (non-renormalizable) terms which are suppressed by inverse powers of Λ and hence will be small, consistent with the success of the standard model.

As a corollary of this argument, one can easily show that the unbroken gauge symmetries of the theory, QCD and electromagnetism, must be vectorlike, i.e., conserve parity, in their coupling to massive fermions for otherwise the quarks and leptons will have left- and right-handed states with different QCD and electromagnetic charges and the argument given above would then require they be massless.

The remarkable correspondence of the states found in the standard model and those expected to remain light in an effective field theory is spoilt when we consider the Higgs sector for no symmetry in the standard model prevents a mass term for the Higgs and the expectation is that the Higgs mass should be of order Λ , the scale of new physics. We have not found a Higgs yet so this might be interpreted as a good result but it also implies that Λ cannot be very high for the Higgs mass sets the scale of electroweak breaking. Moreover, unless the Higgs mass is less than 0 (1 TeV), it is known that perturbative unitarity breaks down¹. Thus we are led to the conclusion that the (successful) concept of the standard model as a low-energy approximation to some underlying theory may be maintained only if the scale Λ of new physics is low; $\Lambda \lesssim 1$ TeV.

What may the new physics be? Three possibilities have been suggested.

(1) The weak interactions may become strong above Λ so that the new physics will be the appearance of non-perturbative phenomena. This possibility and the signals expected were reviewed at last year's Berkeley Conference² so I will not dwell on it, except to comment that, by itself, such a possibility does not answer any of the questions concerning unification of the fundamental interactions and the origin of the many parameters which is at the heart of most theorists' attempts to go beyond the standard model.

(2) There may be composite structure at a scale $\Lambda \approx 1$ TeV. The simplest possibility is that the Higgs are composite and this is the idea pursued in the technicolour theories popular some years ago³. Of course the idea may be extended to include composite fermions or composite gauge bosons even though, as discussed above, symmetry arguments allow them to be light elementary states. Below the composite scale the effective field theory will have non-renormalizable terms $\propto \Lambda^{-n}$ which may give observable effects, e.g., flavour changing neutral currents (FCNC), if Λ is as low as 1 TeV. There may also be structure for the W, Z, and the other states of the theory.

(3) There may be a new symmetry forcing the Higgs mass to be zero up to symmetry breaking effects. The only possible symmetry is supersymmetry and Λ in this case is identified with the supersymmetry breaking scale so the new physics to be expected is the appearance, with masses of $O(\Lambda)$, of the supersymmetric partners of the states of the standard model. Also below the supersymmetry breaking scale we may expect non-renormalizable terms $\propto \Lambda^{-n}$ which may give observable effects such as FCNC.

To determine which, if any, of these possibilities is correct we need to look for signals beyond the standard model. The most obvious signal is the appearance of new states: in composite models there may be technihadrons, subquarks, etc., in supersymmetric or superstring models there should be the supersymmetric partners the gauginos and squarks, etc. An associated expectation is the possibility of new interactions such as new gauge interactions required in versions of the superstring or in composite models or new scalar interactions due to pseudo-Goldstone bosons giving rise to new long range forces. Another source of information about the high energy world is via phenomena which do not occur or are strongly suppressed in the standard model, such as neutrino masses and mixings, nucleon decay, flavour changing neutral currents, etc. Also radiative corrections involving new states may change the standard model predictions - e.g., the ρ parameter, $B^0-\bar{B}^0$ mixing, CP violation - or may give new predictions, e.g., $\sin^2\theta_w$.

It is clear that there are many possible signals of new physics and, of course, none should be neglected. However, in preparing this talk, I have been guided by the (admittedly very few) clouds on the horizon for the standard model in the hope that these are areas where the hope of finding new physics is somewhat enhanced, and also because the existence of such problems allows one to quantify the new physics required. The first cloud is the solar neutrino problem which corresponds to the deficit between the observation of Davis⁴ of 2.1 ± 0.3 SNU (1 SNU = 1 capture/ 10^{36} atoms/sec) and the latest theory estimate from Bahcall et al.⁵ of 8.2 ± 2.5 SNU. The second cloud is the problem of missing matter⁶ - astronomers are broadly in agreement with the observation that our universe seems to be made up mostly of dark (non-luminous) matter - what is it? A third cloud, but of a more controversial nature, is the possible existence of a fifth force; this would certainly signal departure from the standard model and it is interesting to ask how such a force could arise.

In this talk I will try to discuss these problems in the context of recent theoretical ideas about what lies beyond the standard model. In the time available I cannot do justice to more than a small fraction of the theoretical work so to those whose work I am unable to review - my apologies!

2. SUPERSTRING MODELS

2.1. Heterotic string models

By far, most of the theoretical activity of the last year has been on superstrings⁷ and their phenomenological implications. Much of this work has centred on the compactification which is necessary in these models to arrive at a theory with four space-time dimensions.

The original constructions started with the heterotic string⁸, a candidate (finite?) theory unifying the strong, electromagnetic weak and gravitational forces. The compactification proceeds in two stages as illustrated in Fig. 2.1a first constructing an anomaly-free theory in ten space-time dimensions with the 16 remaining dimensions of the left moving states being internal and giving an $E_8 \times E_8$ gauge symmetry group. There is a subsequent stage of compactification to four space-time dimensions⁹, the residual gauge group being some subgroup of $E_8 \times E_8$ - in Calabi-Yau schemes it is $E_6 \times E_6$ - with one supersymmetry generator left unbroken. So far as low energy phenomenology is concerned, it is the E_6 factor which gives the low energy physics - the other E_8 factor, the "hidden" sector of the theory, is thought to be confining at a very high mass scale so all its states are beyond the reach of low energy experiments.

HETEROTIC STRING

FIGURE 2.1a
Calabi-Yau type compactification

Recent work¹⁰ on alternative compactification schemes has shown that it is not necessary to compactify in two steps. This has given rise to what is known colloquially as "4D superstrings", cf. Fig. 2.1b, in which 22 of the 26 dimensions of the left moving states and six of the 10 dimensions of the right moving states are compactified at once leaving just the required four flat dimensions. By comparison with the original construction of Fig. 2.1a, we see that the residual symmetry group formed by this procedure may be much larger, groups up to $SO(44)$ with four supersymmetries being possible¹⁰.

HETEROTIC STRING

FIGURE 2.1b
"4-D" superstring construction

Clearly the profusion of possible compactification schemes is somewhat of an embarrassment for the theorist trying to make "superstring" predictions for low energy phenomenology. It is to be hoped that some criterion to select one out of many will be found but at present we must keep an open mind on the possible structures. Although work towards building realistic models based on the 4D superstrings is in progress, essentially all phenomenological work to date starts from E_6 symmetric models. In view of the above remarks this should be considered as representative of the possibilities but certainly not exhaustive.

2.2. E₆ based models

In models in which the group after the first stage of compactification is E₈×E₈, the low energy group of the visible sector will be a subgroup of E₆. In discussing the possibilities it is easiest to discuss the quantum numbers of states by their SU(3)_C×SU(3)_L×SU(3)_R transformation properties since this group is a subgroup of E₆ and the standard model gauge group SU(3)_C×SU(2)_L×U(1) must be embedded in this [SU(3)]³ structure. In Table 2.1, I give the decomposition of the supermultiplets in the 78 and 27 dimensional representations of E₆

$$\begin{aligned} E_6 &\supset SU(3) \times SU(3) \times SU(3) \\ &\supset SU(3)_C \times SU(2)_L \times U(1) \end{aligned}$$

GAUGE BOSON - GAUGINO SUPERMULTIPLETS

$$\begin{aligned} 78 = & (8,1,1) + (1,8,1) + (1,1,8) \\ & \text{GLUONS } W^{\pm}_{Z,\gamma} + \underline{8 \text{ NEW STATES}} \\ & (+\text{GAUGINOS}) \\ & + (3,3,3) + (\bar{3},\bar{3},\bar{3}) \\ & \underline{J = 1 \text{ LEPTOQUARKS}} \dots \text{MUST BE HEAVY} \\ & (+ \text{INOS}) \end{aligned}$$

FERMION-SFERMION SUPERMULTIPLETS

$$\begin{aligned} 27 &= (1,3,\bar{3}) + (\bar{3},\bar{3},1) + (3,1,3) \\ \text{HIGGS + HIGGSINOS } H_{1,2} & \begin{bmatrix} H_1^0 & H_2^+ & e^+ \\ H_1^- & H_2^0 & \bar{\nu}_e \\ e^- & \nu_R & \underline{N} \end{bmatrix}_R \\ \text{LEPTONS + SLEPTONS } e,\nu,N : & \begin{bmatrix} u \\ d \\ \underline{D} \end{bmatrix}_L + \begin{bmatrix} u \\ d \\ \underline{D} \end{bmatrix}_R \\ \text{QUARKS + SQUARKS } u,d,D & \begin{bmatrix} u \\ d \\ \underline{D} \end{bmatrix}_L + \begin{bmatrix} u \\ d \\ \underline{D} \end{bmatrix}_R \end{aligned}$$

TABLE 2.1

[SU(3)]³ decomposition of 78 dimensional E₆ vector supermultiplet and 27 dimensional E₆ chiral supermultiplet. The underlined states are not in the supersymmetric standard model.

which, according to the superstring picture, must contain the gauge bosons/ gauginos and quarks, leptons, Higgs/squarks, sleptons, shiggs, respectively, of the supersymmetric version of the standard model¹¹. Of course, a specific realization of the compactification process of Fig. 2.1a should give much more detail about the structure of the low energy model - the number of families, the low energy gauge group, the masses and couplings. The simplest possibility is that E_6 should be broken at the compactification scale by Wilson loops¹². In this case one finds¹³ that only 27 (not $\overline{27}$) representations survive and the E_6 group breaks, but no further than $SU(3)_C \times SU(2)_L \times U(1) \times U(1)'$ with an additional Z' gauge boson whose mass should be of order of the electroweak breaking scale. Although a specific Calabi-Yau manifold leading to this group is not known it may plausibly be argued on phenomenological grounds¹⁴ that the multiplet structure at low energies is as given in Fig. 2.2a. This constitutes the "minimal" superstring inspired model and most phenomenological studies have centred on it.

Another possibility, which produces significant changes in the low energy phenomenology, is to allow for the possibility of Higgs in the theory acquiring very large ["intermediate scale" (IS)] vacuum expectation values¹³, thus breaking the gauge group. In the scheme of Fig. 2.2a this is not possible because the absence of $\overline{27}$ representations makes it energetically unfavourable to have large IS breaking but, in the only known Calabi-Yau manifold leading to just three generations¹⁵, it is thought that it is energetically favourable for these to be two stages of IS breaking leading to the model of Fig. 2.2b. Because this model has a known Calabi-Yau manifold many details of its properties are known, including the multiplet structure, the number of generations, and new (discrete) symmetry properties of the low energy Lagrangian which restricts the couplings of the theory. Thus the model deserves to be called a superstring model, illustrating the predictive power to be expected from this approach. Due to the complexity of breaking patterns possible, there is still some uncertainty about the detailed structure of the low energy model¹⁶ but in Ref. 15 it is argued that a pattern of IS breaking is possible leading to the structure of Fig. 2.2b in which the pattern of quark and lepton masses is consistent with experiment. I will use the two examples of Fig. 2.2 as illustrative of the range of possibilities for phenomena following from the superstring.

$$\text{LIGHT STATES} : 3 \times \left[\begin{pmatrix} u \\ d \\ \underline{D} \end{pmatrix}_L + \begin{pmatrix} u \\ d \\ \underline{D} \end{pmatrix}_R + \begin{pmatrix} H_1^0 & H_2^+ & e^+ \\ H_1^- & H_2^0 & \bar{\nu}_e \\ e & \nu_R & \underline{N} \end{pmatrix}_R \right]$$

FIGURE 2.2a

Minimal rank 5 superstring inspired model

2.3. Superstring interactions

The gauge interactions of the low energy effective theory following from the string are fixed by the form of the residual gauge group. The Yukawa and associated scalar interactions are given by some subset of the $SU(3) \times SU(2) \times U(1)$ invariant terms which may be formed using the E_6 multiplets of Table 2.1. The allowed superpotential terms are

$$W = \alpha_1 H_1 Q u^c + \alpha_2 H_2 L \ell^c + \alpha_3 H_2 Q d^c + \alpha_4 H_1 H_2 N + \alpha_5 D D^c N \\ + \lambda_1 D^c L Q + \lambda_2 D \ell^c u^c + \lambda_3 D Y^c d^c + \lambda_4 D Q Q + \lambda_5 D^c u^c d^c \quad (2.1)$$

[The superpotential summarizes the supersymmetry restrictions on couplings. It corresponds to Yukawa couplings in which one identifies any two of the three fields with the fermion component of the superfield and the third the scalar component. It also generates the quartic scalar couplings which are given by $\sum_i |\partial W / \partial \phi_i|^2$.¹¹]

Not all of these terms need be present in any specific compactification scheme for there may be discrete symmetries forbidding some terms. The first terms proportional to α_i in Eq. (2.1) are needed for mass generation, the second set of terms proportional to λ_i may give rise to new effects and the allowed magnitudes of the λ_i have been bounded by a general analysis of these effects¹⁷. Here I will concentrate on some general features.

In models without an intermediate scale there are four possible choices of non-zero λ_i leading to different phenomenological expectations¹⁸ for the baryon number, B , and the lepton number, L , of the new, charge $-1/3$, D states.

$$\begin{aligned} \text{(i)} \quad \lambda_4 = \lambda_5 = 0 & \quad D \text{ is a leptoquark} \quad B(D) = 1/3, L(D) = 1 \\ \text{(ii)} \quad \lambda_1 = \lambda_2 = \lambda_3 = 0 & \quad D \text{ is a diquark} \quad B(D) = -2/3, L(D) = 0 \\ \text{(iii)} \quad \lambda_0 = \lambda_1 = \lambda_2 = \lambda_4 = \lambda_5 = 0 & \quad D \text{ is a quark} \quad B(D) = 2/3, L(D) = 0 \\ \text{(iv)} \quad \lambda_1 = \lambda_2 = \lambda_3 = \lambda_4 = \lambda_5 = 0 & \quad B(D) = 0, L(D) = 0 \end{aligned} \quad (2.2)$$

From this we see how diverse the possibilities are for the nature of the new coloured D state. We will discuss its phenomenological implications in Section 4.2.

In all cases there is a residual R parity¹¹ under which all new SUSY states the gluinos, squarks, sleptons and shiggs are odd. This means they may only be pair produced and the lightest such state (LSP) must be stable¹¹ - in most models it is the photino, $\tilde{\gamma}$. Although R parity may be spontaneously broken¹⁹, the expectation is this will be a small effect.

In models with an intermediate scale of breaking there are many more possibilities because the D quark and ν^c may be heavy (cf. Fig. 2.2b) meaning that the baryon - or lepton - number violating couplings in Eq. (2.1) may be allowed without generating unacceptably fast nucleon decay. For example²⁰ if in Eq. (2.1) only λ_1 and λ_2 are zero, but the ν^c and $\tilde{\gamma}$ are heavy, the remaining terms will not mediate proton decay. In this case there will be mixing between the D_R and d_R states and this leads to the graph of Fig. 2.3 which mediates photino decay into a baryon. This may be a strong enough violation of R parity to invalidate many SUSY searches why rely on the existence of a LSP. It is important, therefore, to look for signals which allow for LSP decay.

Fig. 2.3
Graph leading to photino decay

I will return in Section 4 to a discussion of some of the phenomenological implications of the new interactions given in Eq. (2.1). Their determination would be of great importance for they provide a window on the structure of the underlying multidimensional world which is determined for a specific compactification manifold or scheme.

2.4. Supersymmetric spectrum

The most direct way to verify supersymmetry is to produce and find the new supersymmetric states. Of course, when this is possible depends on the mass spectrum and here I wish briefly to discuss our theoretical expectations for this spectrum and to compare it to the current experimental bounds to find out whether, if supersymmetry is relevant, we should already have found new states.

The most plausible scheme advanced so far for breaking supersymmetry in a superstring model is via a gaugino condensate in the hidden sector²¹. The

supersymmetry breaking will then be communicated to the visible sector via gravitational radiative corrections^{22,23} leading to a gaugino mass as the leading effect. Using this gaugino mass as the dominant source of breaking, further radiative corrections involving the gauge and Yukawa couplings may be reliably computed²⁴ to give the masses of the other supersymmetric states. The resulting spectrum is given in Table 2.2 and applies to any supersymmetry

SUSY PARTICLE	$\tilde{q}$	$\tilde{g}$	$\tilde{\chi}_L$	$\tilde{\chi}_R$	$\tilde{W}$	$\tilde{\gamma}$
	SQUARK	GLUINO	SLEPTONS	WINO	PHOTINO	
MASS (IN UNITS OF $m_{1/2}$)	1.9	1	0.7	0.4	0.3	0.16

TABLE 2.2
Mass expectations for SUSY particles triggered by common
gaugino mass, $m_{1/2}$, at the Planck scale.

breaking scheme in which gauginos acquire the dominant supersymmetry breaking mass in the visible sector. As such it seems to me to be a reasonably reliable estimate of the spectrum in superstring models and I will use it to compare with the experimental bounds. In Table 2.2 there is an unknown overall mass scale, the gaugino mass at the Planck scale, $m_{1/2}$. As a first exercise let us choose $m_{1/2}$ so that the experimental bounds are saturated. From Fig. 2.4 we see that this occurs when the Wino is just at the experimental upper limit of 36 GeV. We may then compare the experimental bounds with the expectation for the other supersymmetric states. From Fig. 2.4 it is clear that the Winos and Zinos are likely to be the first states to be found, followed by the sleptons, with the squarks and gluinos a long way from the current bounds. A more interesting exercise perhaps is to ask what is the most reasonable theoretical expectation for $m_{1/2}$ and compare the resulting spectrum with experiment. This is also done in Fig. 2.4 using for $m_{1/2}$ a value that gives a Higgs mass of $O(300 \text{ GeV})$ without fine tuning of parameters. As such it represents, I think, the most reasonable value. The implication of this may be seen immediately in Fig. 2.4 - we are quite far from producing and observing even the lightest states. Those most likely to be found first are the Wino and Zino if Table 2.1 gives the correct spectrum. The message I wish to draw from this is that the current lack of experimental evidence for supersymmetry is not to be taken as

FIGURE 2.4
Bounds and expectation for SUSY masses. Data from the review of Wu²⁵
See also Refs. 26, 27. The crosses denote the lower bounds obtained
by choosing $m_{1/2}$ in Table 2.1 to give a Wino mass at the current
experimental bound. The dots give the most reasonable expectation for
SUSY masses as detailed in the text.

evidence against supersymmetry - this will not happen until the bounds pass the values corresponding to the dots of Fig. 2.3. Of course it may happen that we are lucky and the SUSY states are lighter than this guesstimate. I hope so!

3. COMPOSITE MODELS

The other main possibility for new physics beyond the standard model is that there should be a composite structure for some, or all, of the states of the standard model. I find this subject more difficult to review than superstring models because there is no accepted dogma - a glance at the submissions on this subject to this conference shows there are as many different models as authors. However a common theme is the generation of new processes and these may be conveniently parametrized by an expansion of the effective Lagrangian in inverse powers of Λ , where here Λ is the scale of compositeness. Thus the new terms are given by

$$\mathcal{L}_{eff} = \frac{\lambda_{1234}^2}{\Lambda^2} (\bar{\psi}_1 \Gamma \psi_2)(\bar{\psi}_3 \Gamma \psi_4) + O\left(\frac{1}{\Lambda^3}\right) \quad (3.1)$$

The four-Fermi interactions should be consistent with the unbroken flavour symmetries and terms allowed by the composite theory will put limits on the theory. The most stringent limits come from the terms in Table 3.1²⁸

FOUR-FERMI OPERATOR	PROCESS	LOWER BOUND ON Λ
$\bar{e} d \bar{u} u^c$	Proton Decay	$\lambda. 10^{13}$ TeV
$\bar{s} d \bar{s} d$	$K^0 - \bar{K}^0$ mixing	$\lambda. 400$ TeV
$\bar{c} u \bar{c} u$	$D^0 - \bar{D}^0$ mixing	$\lambda. 50$ TeV
$\bar{s} d \bar{e} \mu$	$K^+ \rightarrow \mu^+ \mu e^-$ $K_L \rightarrow \mu e^-$	$\lambda. 30$ TeV

TABLE 3.1
Bounds imposed on Λ by various processes

modification in this case is sufficient to suppress FCNC to acceptable rates but that a walking-coupling is not. In view of the approximations involved in computing these graphs, this may be a debatable conclusion but in both cases the message is that the physics of compositeness may be quite different from the physics of QCD making it even harder to determine the phenomenology to be expected from composite models. An interesting footnote on these schemes resurrecting technicolour theories concerns the pseudo-Goldstones of the theory some of which, according to the original theory, should have been very light and should have been found. Because the new schemes enhance the contribution of the technicolour condensates (cf. Fig. 3.1b) the pseudo-Goldstones become more massive, $\gtrsim 0(300 \text{ GeV})$ and will require SSC energies for production³⁴. The only light scalar state is a 0^{++} technidilation^{31,34} - i.e., the resulting low energy theory looks very like the standard model with a 0^{++} Higgs!

3.2. FCNC

Another suggestion²⁸ which leads to the natural suppression of the dangerous processes of Table 3.1 has recently been developed by Georgi and Chivuluka³⁵ for a composite technicolour model in which quarks, leptons and techniparticles are all composite. It relies on mimicing the symmetries of the standard model and, being based on symmetry, may have a wider applicability, than any single composite model. The standard model, without mass terms, has the large $(\text{SU}(3) \times \text{U}(1))^5$ global symmetry of the kinetic terms under which the three flavours in the five independent representations $\begin{pmatrix} u \\ d \end{pmatrix}_L$, u_R , d_R , $\begin{pmatrix} \nu \\ e \end{pmatrix}_L$, e_R transform into each other. Provided the composite model has the same symmetry, there will be no FCNC of the type given in Table 3.1, for these terms are not invariant under this symmetry. Of course, the symmetry must be broken by the mechanism generating quark and lepton masses, m_i , but these effects will give FCNC suppressed by powers of (m_i/Λ) . Thus in the specific realization³⁵ of this general strategy the dangerous $d\bar{s}d\bar{s}$ operator of Table 3.1 arises only via the term $(1/4\pi)^4 \Lambda^{-6} (\phi_L^\mu M_u M_u^\dagger \phi_L)^2$ where M_u is the up quark mass matrices, and is small even for low scales of compositeness.

The advantage of this approach is that one can analyze in some generality the allowed FCNC operators to identify the dominant FCNC effects. This leads to the identification of $K \rightarrow \pi \nu \bar{\nu}$, $B - \bar{B}$ mixing and $e^+e^- \rightarrow \mu^+\mu^-$ non-perturbative structure as the processes most sensitive to the composite structure³⁵.

3.3. Composite Ws

As well as models with composite Higgs, composite quarks and leptons it is possible to construct viable models with composite gauge bosons³⁶. Once again

3.1. Technicolour

Any viable composite model must satisfy these bounds and many have come to grief at this point. For example, a very appealing composite Higgs model was based on technicolour and extended technicolour interactions³, the technicolour force causing a Higgs scalar to result as a bound state of technifermions. However, the model was ruled out²⁹ by too large flavour changing neutral currents (FCNC) mediated by the extended technicolour bosons (see Fig. 3.1a).

FIGURE 3.1a
Quark mass generation and FCNC in extended technicolour models

Further developments³⁰⁻³³ have sought to save the situation arguing that the high energy behaviour of technicolour may be different from colour, perhaps not even asymptotically free^{31,32} or perhaps with a very small β function - "a walking coupling constant"³³. The results for a non-asymptotically free theory may be extracted from the Bethe-Salpeter equation and are given in Fig. 3.1b. The authors of the third entry in Ref. 31 argue that the

FIGURE 3.1b
Quark mass generation and FCNC in extended technicolour models
with a non-trivial fixed point assumed for technicolour.

there are too many possible scenarios to review. The phenomenological implications below the scale of compositeness follow from the generation of structure for the W s and Z such as a large electric dipole moment or anomalous magnetic moment³⁷, and it is clearly of importance to look for any such structure³⁸.

4. PHENOMENOLOGICAL SIGNALS FOR NEW PHYSICS

In this section I will review some of the more interesting signals that have recently been suggested as good indicators of new physics such as supersymmetry, compositeness or some other extension of the standard model.

4.1. Leptoquarks and diquarks

Although these states may occur in both composite and SUSY models, the expectations for their properties is more definite in the latter case, so most phenomenological estimates have been performed for them^{39,40}.

Production may proceed pairwise via gluons or singly via their Yukawa couplings [cf. Eq. (2.1)]. The decay partner differs according to the state. For leptoquarks, model (1) of Eq. (2.2), the decay is via the Yukawa couplings of Eq. (2.1) and so the expectation is that there will be a characteristic signal involving universality violation³⁹ - $\tilde{D} \rightarrow eu = \tilde{D} \rightarrow \mu u$. For diquarks, model (2) of Eq. (2.2), the decay should similarly involve flavour violation $\tilde{D} \rightarrow \bar{u}s, \bar{u}b$.

A point stressed in Ref. 39 is that if one of the $\tilde{D}$ should be lighter than the top the decay mode

$$t \rightarrow e^+ \tilde{D} \quad \text{or} \quad \bar{d} \tilde{D}$$

may be dominant, invalidating the usual signatures for the top quark. Similarly gluino decays may be radically different $\tilde{g} \rightarrow \tilde{D}\bar{D}$. Just as was the case for large R parity violation, the signals usually considered for supersymmetry may not be the correct ones!

4.2. New neutral currents

There have been many recent papers testing for the possible presence of new neutral currents⁴¹⁻⁴³ of the type suggested by the superstring. So far analyses have concentrated on E_6 based models but with the advent of "4-D superstrings" we should allow for the possibility of an (even) more general structure. The group E_6 is of rank 6 so there may be two additional diagonal neutral currents. The most recent and detailed analyses⁴³ use a parametrization which allows for one additional current with vector boson mass eigenstates Z_1^m and Z_2^m given by (in the notation of the first paper of Ref. 43)

$$\begin{aligned} Z_1^m &= Z_1^0 \cos \Theta + Z_2^0 \sin \Theta \\ Z_2^m &= -Z_1^0 \sin \Theta + Z_2^0 \cos \Theta \end{aligned} \quad (4.1)$$

Here Z_1^0 is the standard model Z coupling in the usual way to quarks and leptons and Z_2^0 is the new neutral boson given by

$$Z_2^0 = Z_\psi^0 \cos \beta + Z_\chi^0 \sin \beta \quad (4.2)$$

where Z_ψ^0 and Z_χ^0 are the bosons associated with the U(1) factors in the following breaking pattern

$$E_6 \rightarrow SO(10) \times U(1)_\psi \rightarrow SU(5) \times U(1)_\chi \times U(1)_\psi \quad (4.3)$$

In this parametrization, the "minimal" rank-5 scheme of Fig. 2.2a has $Z_2^0 \equiv Z_\eta$ given by Eq. (4.2) with $\cos \beta = \sqrt{5}/8$. Present data constrains the amount of new neutral current structure (there is no positive evidence for any such structure) and hence gives a lower bound on the mass of the Z_2^m for a given value of β - see Fig. 4.1. Particularly for $Z_2^0 \equiv Z_\eta$ the bound is rather weak

FIGURE 4.1
Bounds on the mass of an additional neutral current
of the type found in E_6 models

($m_{Z_\eta} > 129$ GeV) and it will be important in future to use collider data to improve them (or to find the new Z!). One should perhaps note that the parametrization of Eq. (4.1) is not the most general one for E_6 models for it implicitly assumes the renormalization of different U(1) factors are the same. Relaxing this constraint gives new possibilities and there is even the

possibility that the new current will couple to the combination $(-2Y^L + \sqrt{3} T_3^R + Y^R)$ which decouples from all leptons^{41,42}. In this case the bounds of Fig. 4.1 and Ref. 43 do not apply for the only way such a Z would be visible is via production and decay in hadronic channels, i.e., in the $p\bar{p}$ colliders producing jets⁴⁴

$$p\bar{p} \rightarrow Z_2 + X \quad \begin{array}{c} \nearrow \\ \text{ } \end{array} \begin{array}{c} q\bar{q} \end{array} \quad (4.4)$$

Again the message is that in view of the large number of possibilities, we should look for evidence of new neutral currents in as many channels as possible.

Recent work^{45,46} has also pointed out strong correlations between Higgs structure and neutral current in superstring models. For example, the "minimal" rank-5 scheme of Fig. 2.2a the authors of Ref. 46 find the following bounds on Higgs masses

$$m_{H_2^0} \leq 110-160 \text{ GeV}$$

$$m_{H^\pm, H_3^0} < 0.55 - 1 \text{ TeV} \quad (4.5)$$

These bounds are rather weak at present but improvements on neutral current data may lead to much stronger bounds. Also it should be noted that the minimal rank-5 model has three charged and sixteen neutral Higgs scalars! Indeed in all supersymmetric models the Higgs sector is much richer than in the standard model¹¹ and provides a sensitive testing ground for these schemes.

4.3. Radiative effects

$$\sin^2\theta_W$$

Once again I will concentrate on SUSY effects because only for them is there some universal expectation but clearly compositeness will also make itself felt below the compositeness threshold via radiative corrections.

One of the major successes of Grand Unification was the calculation of $\sin^2\theta_W$ from the unification result together with radiative effects taking care of corrections between the Grand Unified scale and laboratory energies at which $\sin^2\theta_W$ is measured. The experimental measurement of $\sin^2\theta_W$ has so much improved that small differences in radiative corrections are important and, as Prof. Sirlin mentioned in his talk here, there is even an indication that supersymmetric states are needed below the GUT scale to give an acceptable value (first ref. 43)!

$B^0 - \bar{B}^0$ mixing

Radiative corrections probing scales much lower than the GUT scale may also be important and may be used to bound the spectrum of SUSY states^{47,48}. As an example, consider the diagram of Fig. 4.1 which gives a contribution to $B^0 - \bar{B}^0$ mixing. Note that there are expected flavour changing gluino couplings because the matrix U_q diagonalizing the quark mass matrix need not be the same as $\tilde{U}_q$, the matrix diagonalizing the squark mass matrix⁴⁹. Assuming a minimal Higgs structure fixes $\tilde{U}_q$ and the graph of Fig. 4.2 may be calculated. For convenience, I give the result compared to the normal W^+W^- box contribution of the standard model

FIGURE 4.2

Box diagram involving gluinos which gives rise to $\Delta S = 2$ processes. The mixing matrices involved at one vertex are shown.

$$\frac{\text{SUSY BOX}}{\text{NORMAL BOX}} = 3 \left(\frac{\alpha_3}{\alpha_2} \right)^2 \left(\frac{m_W}{m_{\tilde{g}}} \right)^2 \left| \frac{\bar{M}_{\text{SUSY}}}{\bar{M}_{\text{S.M.}}} \right|^2 \quad (4.6)$$

Since the SUSY contribution involves the strong coupling, α_3 , there is a relative enhancement factor of $(\alpha_3/\alpha_2)^2$ which, even for gluino mass $m_{\tilde{g}}$ as large as the W mass, m_W , looks as though it will make the SUSY contribution dominant. However, Bertolini et al.⁴⁸ have recently found a correction to the calculation in which the crossed graph of Fig. 4.1 largely cancels the effect and they find the SUSY box is less than the normal box for $m_{\tilde{g}} = m_{\tilde{b}} \gtrsim 60$ GeV.

One should note that in general one needs a non-minimal Higgs structure to avoid bad fermion mass relations: the resulting change in $\tilde{U}_q$ increases the SUSY contributions^{47,50} so that for supersymmetric masses in the range given above SUSY effects may be significant and may change the standard model predictions for $B - \bar{B}$ mixing. Of course, one would not immediately be able to conclude this was the effect of SUSY, for example a fourth generation would give similar effects⁵¹.

Radiative b decays

In other processes the coupling constant enhancement of gluino contributions may not be cancelled giving even larger radiative corrections. For example, it is found that there is a large $b \rightarrow s + \gamma$. The experimental limit on the branching ratio⁵²

$$\text{B.R. } (B \rightarrow K^* \gamma) < 1.8 \times 10^{-3} \quad (4.7)$$

translates to a bound (assuming $K^* \gamma$ is the dominant decay mode)

$$m_{\tilde{g}} \approx m_{\tilde{q}} > 60-70 \text{ GeV for } m_t \geq 60 \text{ GeV} \quad (4.8)$$

This is comparable to the current bound coming from the CERN $p\bar{p}$ collider^{25,26}, showing how important radiative effects may be. A similar calculation⁴⁸ gives $\text{BR}(b \rightarrow sg) > 10-30\%$ from the SUSY graph with the masses taken as in Eq. (4.8) - approximately 1000 times the non-SUSY contribution! It may well be that B mesons will provide a useful probe of SUSY physics via their non-standard decays.

4.4. Neutrino masses and mixings

Neutrino masses provide a window on the world beyond the standard model for they are forbidden in the standard model - Majorana masses are absent because the standard model is lepton number (L) conserving and Dirac masses do not occur because there is no right-handed neutrino component ν_R . Beyond the standard model there are various neutrino mass generation mechanisms⁵³.

(i) $\Delta L = 2$ Majorana masses. These may occur from an $I = 1$ Higgs, ϕ , - the "Majoron"⁵⁴ - with coupling $\lambda \nu_L \nu_L \phi^{I=1}$ which gives a Majorana mass

$$m_{\nu}^M = \lambda \langle \phi^{I=1} \rangle \quad (4.9)$$

where $\langle \phi^{I=2} \rangle$ is the Higgs vacuum expectation value.

Alternatively they may occur via the Gell-Mann, Ramond, Slansky, Yanagida mechanism via a term $\lambda \bar{\nu}_L \nu_R$ coupling a conventional $I = \frac{1}{2}$ Higgs to the (left-handed) neutrinos to new ν_R states which themselves have a large Majorana mass, M_R . In this case

$$m_{\nu}^M = \frac{\lambda^2 \langle \phi^{I=\frac{1}{2}} \rangle^2}{M_R} \quad (4.10)$$

In (E6 based) superstring theories, the expectation is that the first contribution, Eq. (4.9), is absent because there is no $\phi^{I=1}$ state, but the

second may occur in models with an intermediate scale^{13,15,55} such as that in Fig. 2.2b.

(ii) $\Delta L = 0$ Dirac masses. The natural expectation for a Dirac neutrino mass, if allowed, is of order of the charged lepton mass. It is possible, in the absence of Majorana masses, to arrange a "see-saw" mechanism similar to Eq. (4.10) but involving Dirac masses only⁵⁶. This gives a small Dirac mass to the neutrino but no associated lepton number violation. In superstring models this mechanism may be implemented if there is an intermediate scale⁵⁷. In the minimal rank-5 model, Fig. 2.2a, there is no intermediate scale and no Majorana mass and so Dirac masses must be forbidden by invoking discrete symmetries. However, small Dirac masses may be generated in radiative order^{14,58}.

Considering the origin of neutrino masses, the best theoretical guesses for Majorana mass ratios and mixings may be summarized by:

$$m_{\nu_e} : m_{\nu_\mu} : m_{\nu_e} = m_t^n : m_c^n : m_u^n ; n=1 \text{ or } 2$$

$$(\tau?) \quad (\mu?) \quad (e?)$$

$$\Theta_{ij} \text{ small } \{ \Theta_{\mu\tau} \sim \Theta_{e\mu} \sim \Theta_{\text{CASPERO}}? \quad \Theta_{e\tau} \text{ smaller} \} \quad (4.11)$$

At best these are useful guides to expectations for use in phenomenological estimates.

Consequences of neutrino mass

As Eichler has already reviewed in this conference, Majorana masses must be small to be consistent with the bounds on double β decay and this process also shows no evidence for Majorana emission processes⁵⁹. Non-zero neutrino masses (both Dirac and Majorana) also play a rôle in astrophysics⁶ $\sum_i m_{\nu_i} \lesssim 100$ eV is needed to avoid overclosure of the universe and $\sum_i m_{\nu_i} = 10$ eV would be sufficient for neutrinos to make up the "dark halos". Non-zero masses and mixings could also account for the solar neutrino deficit by the oscillation of the electron neutrinos into neutrinos of a different type. For oscillation in vacua large mixing angles are needed with $\Delta m_{ij}^2 \equiv m_{\nu_i}^2 - m_{\nu_j}^2 > 10^{-12} (\text{eV})^2$ outside the theoretically favoured range of Eq. (4.11). However, oscillation in the presence of matter is enhanced by the interaction - the MSW effect⁶⁰ - and can give effects sufficient to explain the solar neutrino deficit for small mixing angles consistent with Eq. (4.11). For example⁶¹, $\Delta m_{e\mu}^2 = 10^{-4} - 10^{-7} (\text{eV})^2$ and $|U_{e\mu}|^2 > 10^{-3}$ gives sufficient $e \rightarrow \mu$ oscillation, or $\Delta m_{e\mu}^2 = 10^{-8} - 10^{-4} (\text{eV})^2$ and $|U_{e\mu}|^2 > 10^{-4}$ gives sufficient $e \rightarrow \mu$ mixing. In the former case, using the theory expectations of Eq. (4.11), one finds $\Delta m_{e\mu}^2 = 10^{-4} - 10^4 (\text{eV})^2$, good news for lab searches.

The supernova events have also been used to bound neutrino masses, but it seems that a conservative estimate gives bounds somewhat less than current laboratory bounds⁶². If some of the events of Kamiokande are ν_e for neutronisation one gets further bounds on Δm^2 and mixing angles from the requirement that the MSW effect should not have taken place. This rules out most of the region which would give sufficient oscillation to solve the solar neutrino deficit - only the $e \rightarrow \mu$ non-adiabatic solution survives⁶³. A cunning way out of this⁶⁴ is if the parameters are just right to rotate the neutrinos from the supernova back to ν_e^s as they pass through the earth! However, all these conclusions should be treated with caution as the evidence that some of the Kamiokande events are ν_e^s is not very strong⁶².

Neutrino magnetic moment

Another explanation of the solar neutrino depletion relies on it having a magnetic moment⁶⁵, for then there will be a rotation of the left-handed neutrino ν_{λ_L} into a right-handed ν_{λ_R} component which has no weak interactions in the laboratory experiments. Such an explanation would suggest a correlation with sunspot activity but the evidence for this is tenuous at present. The experimental bound on a neutrino magnetic moment is⁶⁶

$$\mu < 1.5 \times 10^{-10} \mu_B \quad \left(\mu_B = \frac{e}{2m_e} \right) \quad (4.12)$$

What μ is needed to explain the solar neutrino deficit? The energy of a neutrino with $\mu = 10^{-10} \mu_B$ in a magnetic field of 10^3 gauss typical of the surface of the sun is $(\mu B)_0 \equiv 3 \cdot 10^{-11} \text{ cm}^{-1}$. On the other hand, ν_{λ_L} has coherent interaction with matter giving it, in the convective zone of the sun, a potential energy $V \approx 30 (\mu B)_0$ inhibiting the $\nu_{\lambda_L} \rightarrow \nu_{\lambda_R}$ rotation. This leads to an estimate⁶⁷ for the necessary μB to explain the solar neutrino depletion of

$$\mu B \geq (6-12) (\mu B)_0 \quad (4.13)$$

This suggests that a large $\mu \sim 10^{-10} \mu_B$ is needed, close to the experimental lower bound of Eq. (4.12). How could such a large moment arise? In the standard model it is very small⁶⁸. One might hope that new physics at a scale Λ^2 could give big effects generating a term of the $SU(3) \times SU(2) \times U(1)$ invariant form⁶⁷

$$\frac{1}{\Lambda^2} \lambda_2 (\bar{\nu}_L \bar{e}_L) \sigma_{\mu\nu} \nu_R \begin{pmatrix} \phi^0 \\ \phi^- \end{pmatrix} B^{\mu\nu} \quad (4.14)$$

where (ϕ_-^0) is the Higgs doublet and $B^{\mu\nu}$ the U(1) field strength. The problem with this is that the term of Eq. (4.14) gives rise in radiative order to a neutrino mass by absorbing the B field on the ν_L componed. Naively this gives $\delta m_\nu \sim g/4\pi^2 \mu \Lambda^2$ requiring $\Lambda < 1$ GeV, surely too low. Of course, this argument may be false because m_ν must be renormalized (but why is the result so small?). If one ignores this problem then one may build models to generate a term like Eq. (4.14). For example, in Ref. 69, acceptably large values of μ are generated in a model involving new scalar fields of mass $< 10^3$ GeV. I think the lesson to take from this is that to get a sufficiently large magnetic moment there must be new physics at a low scale.

A new massive neutrino (wimps, cosmions, magninos...)

Here I wish to talk briefly about another possibility for solving the solar neutrino problem which may explain the missing (dark) matter at the same time. The explanation requires the existence of a new neutral fermion which must be a weakly interacting and massive particle (hence the name WIMP). Its properties, following Gilliland et al.⁷⁰, may be summarized as follows:

- (i) a low annihilation rate in the sun,
- (ii) a neutral fermion with mass M in the range $4 < M < 10$ GeV,
- (iii) a large cross-section for scattering within the sun. $\sigma \sim 4 \cdot 10^{-36}$ cm² to be compared with a typical weak interaction cross-section of $4 \cdot 10^{-38}$ cm².

The latter property is required for the efficient transport of energy from the centre of the sun - a 10% reduction in its temperature leads to a 2/3 reduction in the neutrino emission as required to explain the observed solar neutrino flux. The mass bounds result from the requirement that the inner core of the sun is cooled without affecting the outer luminous layers and thus leaving the luminosity essentially unaffected.

There are many candidates for a new neutral lepton; it could be a fourth generation lepton, it could be the ν_R or N state of Table 2.1 following from the E_6 multiplet structure of the superstring inspired models... The various models can best be characterized by the way in which a large cross-section is achieved:

- (a) via an anomalously large magnetic moment (the magnino⁷¹),
- (b) via a large neutral current interaction⁷²,
- (c) in a large scalar interaction⁷³.

- (a) The magnino solution starts with a fourth generation of leptons, m , with a Dirac mass and a conserved particle number. As a result, they will not annihilate in the sun if there is a particle-antiparticle asymmetry in this lepton.

It is possible, for the mass in the allowed mass range (i), that the magnino may be a component of the dark matter, thus combining the solution to the two astrophysical problems.

A large magnetic moment arises in this model via the graph of Fig. 4.3, where h^+ is a member of a new Higgs doublet. To get a large enough magnetic moment to satisfy condition (ii) above, the Higgs coupling must be of order 1 and $m_{h^-} \approx m_{\tilde{m}^-} \approx m_{m^0}$. The model has thus many new light states and hence many experimental implications. The obvious signal is:

FIGURE 4.3
Graph generating a large magnino magnetic moment

$$e^+e^- \rightarrow m^+m^- \rightarrow m^0 + e^+ + \gamma + \bar{m}^0 + e^- + \bar{\gamma} \quad (4.15)$$

and in order to explain why this is not seen for masses given by (ii) it is necessary to have m^+ and m^0 very close in mass⁷¹. This window is almost closed by the recent results of Perl et al. and the TPC/2 γ ⁷⁴ experiments which leave only the range $(m_{\tilde{m}^-} - m_{m^0}) < 0.4$ GeV as a possibility. Because the neutrino is light it will contribute of the order of 14% to the Z width. The neutral Higgs should also be visible via

$$e^+e^- \rightarrow h^0 \bar{h}^0 \rightarrow 2\gamma \quad (4.16)$$

and the non-observation of this⁷⁵ gives the bound⁷¹ $m_{h^0} > 12.7$ GeV suggesting $m_{h^0} > m_{h^+}$ in which case h^+ should decay via $h^+ \rightarrow \nu e^+ \gamma$.

(b) The possibility that the heavy lepton has an anomalously large cross-section as a result of a new neutral current interaction seems difficult to achieve because the cross-section required needs $m_{Z'} < m_Z$ for reasonable values of the coupling. The only possibility is the Z' decouples from light leptons; in this case it would only be observed in hadronic colliders $p\bar{p} \rightarrow Z' \rightarrow$ hadronic jets. Although the CERN collider experiments²⁵ give limits on this, it is possible to obey these limits while still giving the WIMP a large enough

cross-section for a Z' mass $\gtrsim 0(60 \text{ GeV})^{72}$. As discussed in Section 4.2, this scheme may be implemented in a superstring model with z' a combination of E_6 currents and the new neutral lepton a member of the 27-dimensional representation of E_6 .

(c) Somewhat surprisingly it is possible to give a large enough cross-section with the standard model Higgs of mass 1.7 GeV^{76} . This mass range has only recently been ruled out by the CUSB and CLEO results announced at this conference⁷⁷. Two doublet Higgs, as required in supersymmetric models, have more flexibility due to the two vacuum expectation values but again it seems the recent limits on neutral Higgs can rule them out. Non-standard scalar interactions will, of course, be able to generate a large enough cross-section with $M_{\phi 0} > 5 \text{ GeV}^{73}$.

Once again, I think the encouraging moral to be drawn is that attempts to explain the solar neutrino problem lead to new physics which is accessible to present experiments.

5. THE FIFTH FORCE ?

5.1. Experimental situation

I turn lastly to a brief discussion of the developments concerning the fifth force. The situation at last year's Berkeley conference may conveniently be summarized by writing the total potential between two masses m_a and m_b in the form

$$V(r) = -G_{\infty} \frac{m_a m_b}{r} \left(1 + \alpha_{ab} e^{-r/\lambda} \right) \quad (5.1)$$

The primary evidence for the claimed fifth force comes from geophysical experiments and from experiments of the Eötvös, Pekar and Febete type. The former experiments are broadly independent of the chemical compositions of the interacting bodies and favour a repulsive force with⁷⁸

$$\alpha_{ab} = -\xi \approx -(1.0 \pm 0.4) 10^{-2} \quad (5.2)$$

The latter experiments suggest a dependence on chemical composition, the favoured interpretation being that the force couples to baryon number B , and gives⁷⁸

$$\alpha_{ab} = - \left(\frac{B_a}{\mu_a} \right) \left(\frac{B_b}{\mu_b} \right) \xi$$

$$\xi \approx 10^{-2}$$

$$1 \text{ m} < \lambda < 10^4 \text{ m} \quad (5.3)$$

Since the Berkeley conference two more results have been published. Thieberger⁷⁹ using a hollow copper tube floating in water near a cliff finds a non-null result,

$$\xi \lambda \approx (1.2 \pm 0.4) \text{ m} \quad (5.4)$$

whereas Stubbs et al⁸⁰, using a torsion balance containing Cu and Be test bodies find no effect and give the bound

$$|\xi \lambda| \leq 0.4 \text{ m} \quad (5.5)$$

both results being for a nominal value $\lambda = 100 \text{ m}$. These new results are in contradiction (ignoring possible systematic errors) unless Eq. (5.3) is modified to allow for a force coupling to a current other than baryon number. For example, if B in Eq. (5.3) is replaced by

$$q = B \cos \theta_5 + L \sin \theta_5 \quad (5.6)$$

where L is the lepton number, then a specific choice of θ_5 could selectively reduce the effect for the Cu-Be experiment. In particular if $\theta_5 = -11^\circ$ the anomalous effect vanishes for Cu-Be but is still appreciable for Cu-H₂O. To test this possibility, Adelberger et al. repeated their experiment using Al and Be test bodies and again found a null result⁸¹. They find that only for a narrow range with $\theta_5 \approx -63^\circ$ can the two experiments be reconciled and in this case the terrestrial effect should vanish. The only way to reconcile all experiments is to allow for more than one new Yukawa interaction in Eq. (5.1) giving

$$V(r) = -G_N \frac{m_a m_b}{r} \left(1 \pm \sum_{\ell=1}^N q_{a\ell} q_{b\ell} \beta_\ell e^{-r/\lambda_\ell} \right) \quad (5.7)$$

where a positive sign corresponds to a scalar or tensor exchange and negative sign to a vector exchange. In view of the uncertainties associated with the various experiments⁸⁴ I will not pursue this analysis further but turn to the general question of what could be the origin of new long-range forces.

5.2. Possible origins for long range forces

The geophysical evidence favours a repulsive force which requires a vector exchange. This was discussed by Bars and Visser⁸² who argued that very weak vector forces may emerge from Kaluza-Klein compactification. So far as I am aware there has not been significant development of this idea as it proves to be difficult to build realistic models without fine tuning to keep the interaction strength sufficiently weak.

A more standard source of weak long-range forces comes from the possible existence of pseudo-Goldstone bosons. In their usual guise these have pure pseudoscalar couplings and do not give rise to coherent interactions but it is possible for them to have a scalar component in their coupling too. Peccei, Sola and Wetterich have developed this idea under the assumption that the pseudo-Goldstone boson is the dilaton or "cosmon" coupled to the divergence of the dilation current. This current is anomalous due to the mass scale that enters on renormalization. As a result the would-be Goldstone mode acquires a mass, m_s , given by

$$m_s^2 = \frac{m^4}{M^2} \quad (5.8)$$

where m is a typical scale generated by the dilatation anomalies and M is the scale of the spontaneous breakdown of the dilatation symmetry which give rise to the Goldstone mode. Peccei et al. argue that $m > \Lambda_{\text{QCD}}$, where Λ_{QCD} is the scale introduced by QCD, and, choosing for M the Planck mass M_p they find

$$m_s \geq 10^{-11} \text{ eV} \quad \text{i.e.} \quad \lambda \leq 10^4 m \quad (5.9)$$

The force resulting from the interchange of the cosmon between two nuclei is

$$V_{N,N'} = - \frac{G_N}{4\pi} \frac{Q_N Q_{N'}}{r} e^{-r/\lambda} \quad (5.10)$$

where

$$Q_N = f \langle N | \theta^\mu_\mu | N \rangle \approx f \left(1 - \frac{\beta}{40}\right) \quad (5.11)$$

and θ^μ_μ is the anomalous trace of the energy momentum tensor breaking the dilatation symmetry and f an effective coupling.

Note that the force given by Eq. (5.10) is attractive and cannot give the geophysical result Eq. (5.2). Following Ref. 84, the authors of Ref. 83 argue that this experiment gives only a 2σ non-null effect and choose to interpret

the result as placing an upper bound $|\alpha| < 10^{-2}$ for λ in the range between a few metres and a kilometer. The point out that the B dependence is, however, repulsive since

$$Q_i Q_j = f^2 \left(1 - \frac{B}{2\alpha} \dots \right) \quad (5.12)$$

From the geophysical bounds on the leading term they find

$$f^2 \lesssim 10^{-2} \quad (5.13)$$

and so the composition dependent term has magnitude

$$\frac{f^2}{2\alpha} \lesssim 5 \cdot 10^{-4} \quad (5.14)$$

to be compared with the reported signal of Ref. 79 of $\xi = 3 \cdot 10^{-3}$ for $\lambda \approx 10^3$ m. I think this exercise illustrates quite nicely that a pseudo-Goldstone boson with a reasonable set of parameters can give rise to measurable (attractive) weak long-range forces in the range given by Eq. (5.3). However, an argument against this optimistic result is the fact that a pseudo-Goldstone particle with mass and coupling given by Eqs. (5.9) and (5.11) respectively leads to an unacceptable cosmology⁸⁵. One should perhaps still follow the advice of the first author of Ref. 83 who last year stated that it was "premature to fit the fifth force"³⁶.

6. CONCLUSIONS

Clearly it is impossible to summarize briefly the ideas for physics beyond the standard model - they are just too numerous. However, a guiding theme has been the problem of mass generation and nearly all approaches try to shed light on this subject and to explain the origin of the electroweak breaking mass scale and the fermion masses. My own prejudice is that the supersymmetric explanations of the origin of the mass hierarchy are the most economical and persuasive. I am not (yet) discouraged by the lack of any substantial evidence for departure from the standard model for I think that the most reasonable expectation for the masses of the new supersymmetric states puts them currently out of reach of experiment (cf. Fig. 2.4) and that we have to wait until the new machines allow us to cross the threshold for production of the new states. Once that happens it should be an exciting time!

7. DISCUSSION

Question: B.F.L. Ward (Tennessee): Your estimate of what were reasonable sparticle masses seems to be on the high side, in view of the existence of models which are perfectly consistent and which allow smaller sparticle masses. Did you have any reason for eliminating the latter models?

Answer: G.G. Ross (Oxford): The reason I believe my estimate to be reasonable is that the Higgs mass, which sets the electroweak scale, has a mass typical of a lepton. The leptons in the table are of the order of 100-200 GeV, and the electroweak scale is of that order. But I take your point. We may be lucky and supersymmetric states may be lighter than this estimate. I hope so.

Comment: R.E. Marshak (Virginia Polytech.): The speaker has spent most of his time on E_6 -inspired phenomenology beyond the standard model, which is his privilege. However, since it is not even clear that superstring theory predicts correctly the phenomenology of the standard model in the low energy limit, I should like to reiterate the salient features of $SO(10)$ -inspired phenomenology beyond the standard model which, like E_6 , abhors the "desert" of $SU(5)$ (already excluded by the failure to observe proton decay). The physical basis for $SO(10)$ grand unification path is the emergence of $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ as the maximum gauged anomaly-free chiral flavour group within the framework of QCD for an arbitrary number of flavours (see my contributed paper to this conference). The new phenomenology that follows from this simple observation includes (finite mass) Majorana neutrinos, a possible second low-mass neutral weak boson, an electric dipole moment of the neutron at the level of 10^{-26} e cm, detectable $K_L \rightarrow \mu \bar{e}$ decay and neutron oscillations (it should be noted that E_6 does not allow for detectable neutrino oscillations). Clearly, the much more conservative $SO(10)$ -inspired phenomenology beyond the standard model is already quite rich.

Answer: G.G. Ross: I did not want to imply that E_6 was the unique possibility. I think there are lots of interesting structures. My brief was to talk about phenomenology, and it happens that the phenomenology that has been explored since the last conference has been mostly E_6 . I certainly do not advocate that you should restrict yourself just to E_6 .

Question: N. Dombey (Sussex): It is very important to verify that the self-interactions of the gauge bosons are given experimentally by the values predicted in the standard model to test the local gauge group $SU(2) \times U(1)$. This will not be done until $e^+e^- \rightarrow W^+W^-$ is measured.

Answer: G.G. Ross: It is indeed valuable to try to identify the most sensitive tests of the standard model.

REFERENCES

- 1) B.W. Lee, C. Quigg and H.B. Thacker, Phys. Rev. D16 (1977) 1519.
- 2) M.S. Chanowitz, Proceedings of the XXIII International Conference on High Energy Physics, ed. S. Loken (World Scientific, Singapore, 1986) 445.
See also: R. Casalbuoni, S. De Curtis, D. Dominici and R. Gatto, Nucl. Phys. B282 (1987) 235 and references therein.
- 3) For a review see: E. Farhi and L. Susskind, Physics Reports 74 (1981) 277.
- 4) For a review see: Solar Neutrinos and Solar Astronomy, American Institute of Physics, eds. M.L. Cherry, W.A. Fowler and K. Lande (1985).
- 5) J.N. Bahcall, W.F. Huebner, S.H. Lubow, P.D. Parker and R.K. Ulrich, Rev. Mod. Phys. 54 (1982) 767.
The latest result from Bahcall et al. was communicated to me by S.P. Rosen.
See: S.P. Rosen and J.M. Gelb, Harrogate Conference, England, August 1986.
- 6) For a recent review, see: J. Silk, Proceedings of the XXIII International Conference on High Energy Physics, ed. S. Loken (World Scientific, Singapore, 1986) 465.
- 7) M. Green and J. Schwarz, Phys. Lett. 149B (1984) 117.
- 8) D.J. Gross, J.A. Harvey, E. Martinec and R. Rohm, Phys. Rev. Lett. 54 (1985) 502; Nucl. Phys. B260 (1985) 569.
- 9) P. Candelas, G. Horowitz, A. Strominger and E. Witten, Nucl. Phys. B258 (1985) 46.
- 10) K.S. Narain, Phys. Lett. 169B (1986) 41;
K.S. Narain, M.H. Sarmadi and E. Witten, Nucl. Phys. B279 (1987) 369;
K.S. Narain, M.H. Sarmadi and C. Vafa, Harvard Preprint HUTP-86/A089 (1986);
H. Kawai, D.C. Lewellen and S.H.H. Tye, Phys. Rev. Lett. 57 (1986) 1832; Nucl. Phys. B288 (1987) 1;
W. Lerche, D. Lust and A.N. Schellekens, Nucl. Phys. B287 (1987) 477;
M. Müller and E. Witten, Phys. Lett. 182B (1986) 28;
I. Antoniadis, C. Bachas and C. Kounnas, Nucl. Phys. B289 (1987) 87;
A.N. Schellekens, CERN Preprint TH. 4807 (1987);
L.E. Ibañez, H.P. Nilles and F. Quevedo, Phys. Lett. 187B (1987) 25;
L.E. Ibañez, J.E. Kim, H.P. Nilles and F. Quevedo, Phys. Lett. 191B (1987) 3;
A. Chamseddine and J.-P. Derendinger, Preprint ETH-PT/87-2 (1987).
- 11) For a review of the supersymmetric standard model, see: H.P. Nilles, Physics Reports 110 (1984) 1;
H. Haber and G.L. Kane, Physics Reports 117 (1983) 75.
- 12) Y. Hosotani, Phys. Lett. 126B (1983) 309; 129B (1983) 193;
E. Witten, Phys. Lett. 126B (1984) 351.
- 13) M. Dine, V. Kaplunovsky, M. Mangano, C. Nappi and N. Seiberg, Nucl. Phys. B259 (1985) 519;
J.D. Breit, B.A. Ovrut and G.C. Segrè, Phys. Lett. 158B (1985) 33;
F. del Aguila, G. Blair, M. Daniel and G.G. Ross, Nucl. Phys. B2872 (1986) 413.

- 14) B.A. Campbell, J. Ellis and D.V. Nanopoulos, Contributed paper Nr. 104 and CERN Preprint TH. 4504 (1986);
E. Cohen, J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys. Lett. 165B (1985) 76;
J. Ellis, K. Enqvist, D.V. Nanopoulos and F. Zwirner, Mod. Phys. Lett. A1 (1986) 57; Nucl. Phys. B276 (1986) 14.
- 15) S.-T. Yan, in Argonne Symposium on Anomalies, Geometry and Topology (World Scientific, 1985);
B.R. Greene, K.H. Kirklin, P.J. Miron and G.G. Ross, Nucl. Phys. B278 (1986) 667; B292 (1987) 606; Phys. Lett. B180 (1986) 69.
- 16) J. Ellis, K. Enqvist, D.V. Nanopoulos and K.A. Olive, CERN Preprint TH. 4713 (1987);
S. Kalara and R.N. Mohapatra, University of Maryland Preprint UM-TH-598-87 (1987).
- 17) B.A. Campbell et al., Contributed paper Nr. 105 and CERN Preprint TH. 4473 (1986).
- 18) V. Barger, N.G. Deshpande and K. Hagiwara, Contributed paper Nr. 251 and Preprint MAD/PH/347 (1985) 375;
- 19) G.G. Ross and J.W.F. Valle, Phys. Lett. 151B (1985) 375;
J. Ellis, G. Gelmini, C. Jarlskog, G.G. Ross and J.W.F. Valle, Phys. Lett. 150B (1985) 142.
- 20) M.C. Bento, L. Hall and G.G. Ross, Nucl. Phys. B292 (1987) 400.
- 21) J.-P. Derendinger, L.E. Ibañez and H.P. Nilles, Phys. Lett. 155B (1985) 65;
M. Dine, R. Rohm, N. Seiberg and E. Witten, Phys. Lett. 156B (1985) 55.
- 22) J.-P. Derendinger, L.E. Ibañez and H.P. Nilles, Nucl. Phys. B267 (1986) 365;
K. Choi and J.E. Kim, Phys. Lett. 165B (1985) 71;
L.E. Ibañez and H.P. Nilles, Phys. Lett. 169B (1986) 354.
- 23) J. Ellis, D.V. Nanopoulos, M. Quiros and F. Zwirner, Phys. Lett. 180B (1986) 83;
J. Ellis, A.B. Lahanas, D.V. Nanopoulos, M. Quiros and F. Zwirner, Phys. Lett. 188B (1987) 408;
P. Binétruy, S. Dawson and I. Hinchliffe, Phys. Lett. 179B (1986) 262 and Phys. Rev. D35 (1987) 2215;
N. Dragon, U. Ellwanger and M.G. Schmidt, Heidelberg Preprint HD-THEP-86-21 (1986).
- 24) L.E. Ibañez and G.G. Ross, Phys. Lett. B110 (1982) 45;
L. Alvarez-Gaumé, M. Claudson and M.B. Wise, Nucl. Phys. B207 (1982) 96;
C.R. Nappi and B.A. Ovrut, Phys. Lett. B113 (1982) 175;
K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, Progr. Theor. Phys. 68 (1982) 927;
M. Dine and W. Fischler, Phys. Lett. B110 (1982) 227;
J. Ellis, K. Enqvist, D.V. Nanopoulos and F. Zwirner,
- 25) T.T. Wu, Rapporteur talk at this Conference.
- 26) S. Greer, Review talk at EPS High Energy Physics Conference, Uppsala, Sweden (1987).

- 27) H.J. Behrend et al., Contributed paper Nr. 206 and DESY Preprint 87-013 (1987);
M. Artuso et al., Contributed paper Nr. 289.
- 28) For a recent review see: I. Bars, International Conference on Superstrings, Cosmology and Compositeness, Maryland (1987) and Preprint USC-87/HEP-2.
- 29) S. Dimopoulos, S.A. Raby and P. Sikivie, Nucl. Phys. B176 (1980) 449;
S. Dimopoulos and J. Ellis, Nucl. Phys. B182 (1981) 505; Nucl. Phys. B193 (1981) 150;
E. Eichten and K.D. Lane, Phys. Lett. 90B (1979) 237.
- 30) B. Holdom, Phys. Rev. D24 (1981) 1441.
- 31) K. Yamawaki, M. Bando and K. Matumoto, Phys. Rev. Lett. 56 (1986) 1335;
K. Yamawaki, Contributed paper Nr. 292 and Preprint DPNU-87-18 (1987);
M. Bando, T. Morozumi, H. So and K. Yamawaki, Contributed paper Nr. 125 and Preprint DPNU-87-10 (1987).
- 32) W.A. Bardeen, C.L. Leung and S.T. Love, Phys. Rev. Lett. 46 (1986) 1230;
C.L. Leung, S.T. Love and W.A. Bardeen, Nucl. Phys. B273 (1986) 649.
- 33) T. Applequist, D. Karabali and L.C.R. Wijewardhana, Phys. Rev. Lett. 57 (1986) 957;
T. Applequist and L.C.R. Wijewardhana, Phys. Rev. D35 (1987) 774.
- 34) W.-C. Kuo, B.-L. Young and D. McKay, Contributed paper Nr. 199.
- 35) R.S. Chivuluka and H. Georgi, Harvard University Preprints HUTP-87/A003, HUTP-87/A036 or BUHEP-87-14 (1987).
- 36) For a review and references, see: R. Peccei, Proceedings of the XXIII International Conference on High Energy Physics, ed. S. Loken (World Scientific, Singapore, 1986), 3.
- 37) For recent work on this subject, see: M.A. Samuel and T.G. Rizzo, Contributed paper Nr. 7 and Phys. Rev. D35 (1987) 403;
G. Couture, J.N. Ng, J.L. Hewett and T.G. Rizzo, Contributed paper Nr. 46 and Preprint DOE-ER40200-089 (1987);
F. Budjema and N. Dombey, University of Sussex Preprint (1987).
- 38) J. Maalampi, D. Schildknecht and K.-H. Schwartz, Phys. Lett. B166 (1986) 361;
M. Kuroda, J. Maalampi, K.-H. Schwartz and D. Schildknecht, Nucl. Phys. B284 (1987) 271; Phys. Lett. 190B (1987) 217.
- 39) V. Barger, N.G. Deshpande and K. Hagiwara, Contributed paper Nr. 251 and Preprint MAD/PH/347 (1987).
- 40) J. Gunion and E. Ma, Contributed paper Nr. 200 and Preprint UCD-87-08 (1987);
F. del Aguila, E. Laermann and P. Zerwas, University of Florida Preprint UFTP-87-11 (1987);
J.L. Hewett and S. Pakvasa, University of Hawaii Preprint UH-511-624-87 (1987).
- 41) S.M. Barr, Phys. Rev. Lett. 55 (1985) 2778;
E. Cohen, J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys. Lett. 165B (1985) 76;

- V. Barger, N.G. Deshpande and K. Whisnant, Phys. Rev. Lett. 56 (1986) 30;
 L.S. Durkin and P. Langacker, Phys. Lett. 166B (1986) 436;
 F. del Aguila, G. Blair, M. Daniel and G.G. Ross, Nucl. Phys. B283 (1987) 50;
 D. London and J.L. Rosner, Phys. Rev. D34 (1986) 1530.
- 42) F. del Aguila, Ll. Ametller, R.D. Field and Ll. Garrido, University of Florida Preprint UAB-FT-175 (19);
 F. del Aguila, J.A. Gonzalez and M. Quiros, University of Michigan Preprint UM-TH-876-11 (1987);
 V. Barger and K. Whisnant, Contributed papers Nr. 180 and 57 and Preprints MAD/PH/351, MAD/PH/336 (1987);
 T.G. Rizzo, Ames Laboratory Preprint IS-J 2280 (1987).
- 43) U. Amaldi, A. Bohm, L. Durkin, P. Langacker, A. Mann, W. Marciano, A. Sirlin and H. Williams, Contributed paper Nr. 254 and CERN Preprint EP/87-93 (1987);
 G. Costa, J. Ellis, G.L. Fogli, D.V. Nanopoulos and F. Zwirner, Contributed paper Nr. 290 and CERN Preprint TH. 4675 (1987).
- 44) UA2 Collaboration, Phys. Lett. 194B (1987) 158.
- 45) P. Roy, Proceedings of the XXIII International Conference on High Energy Physics, ed. S. Loken (World Scientific, Singapore, 1986) 387;
 H.E. Haber and M. Sher, Preprint SCIPP 86/66 (1986);
 J. Ellis, D.V. Nanopoulos, S.T. Petcov and F. Zwirner, CERN Preprint TH. 4454 (1986);
 M. Drees, Preprint MAD/PH/313 (1986).
- 46) J.F. Gunion, L. Roszkowski and H.E. Haber, Preprint UCD-86-41 (1986).
- 47) L.J. Hall, V.A. Kostelecky and S. Raby, Nucl. Phys. and references therein.
- 48) S. Bertolini, F. Borzumati and A. Masiero, Submitted papers Nr. 145, 77 and Preprints CMU-HEP86-20, CMU-HEP86-19 (1986).
- 49) A. Bouquet, J. Kablan and C.A. Savoy, Preprint LPTHE 8425 (1983);
 J. Donaghue, H.P. Nilles and D. Wyler, Phys. Lett. 128B (1983) 55;
 M.J. Duncan, Nucl. Phys. B221 (1983) 285;
 M.J. Duncan and J. Trampetic, Phys. Lett. 134B (1984) 439;
 R. Barbieri and R. Gatto, Phys. Lett. B110 (1982) 210.
- 50) G.G. Ross and S. Wilkinson, Nucl. Phys. B282 (1987) 382.
- 51) A. Datta, E.A. Paschos and W. Turke, University of Dortmund Preprint DO-TH 87/9 (1987).
- 52) P. Avery et al., CLEO Collaboration Preprint CLNS 86/737 (1986).
- 53) For a recent review on neutrino masses, see: H. Harari and Y. Nir, Submitted papers Nr. 223 and 224 and Preprints SLAC-PUB-4137, 4224 (1987).
- 54) G. Gelmini and M. Roncadelli, Phys. Lett. 99B (1981) 411.
- 55) S. Nandi and U. Sarkar, Phys. Rev. Lett. 56 (1986) 564.

- 56) S. Bertolini and J. Liu, Contributed paper Nr. 184 and Preprint CMU-HEP-87-09 (1987);
L. Wolfenstein and D. Wyler, Nucl. Phys. B218 (1983) 205;
M. Roncadelli and D. Wyler, Phys. Lett. 133B (1983) 325;
P. Roy and O. Shanker, Phys. Rev. Lett. 52 (1984) 713.
- 57) R.N. Mohapatra, Phys. Rev. Lett. 56 (1986) 561;
R.N. Mohapatra and J.W.F. Valle, Phys. Rev. D34 (1986) 1642;
J.L. Rosner, Comments in Nucl. and Part. Phys. 15 (1986) 195;
D. London, G. Belanger and J.N. Ng, Phys. Rev. D34 (1986) 2867.
- 58) K. Olive and D.V. Nanopoulos, Madison preprint MAD-TH 86-27 (1986).
- 59) D.O. Caldwell et al., Submitted paper Nr. 120 (1987);
P. Fisher et al., Preprint CALT-63-487 (1987).
- 60) S.P. Mikheyev and A.Yu. Smirnov, Sov. J. Nucl. Phys. 42 (6) (1985) 913;
Nuovo Cimento 9C (1986) 17;
L. Wolfenstein, Phys. Rev. D17 (1978) 2369.
- 61) See: S.P. Rosen and J.M. Gelb, Ref. 5), for a recent review.
- 62) D. Schramm, Rapporteur talk at this conference.
- 63) T.K. Kuo and J. Pantaleone, Contributed paper Nr. 149 and Preprint PURD-TH-87-08 (1987).
- 64) J. Arafune, M. Fukugita, T. Yanagida and M. Yoshimura, Contributed paper Nr. 195 and Kyoto University Preprint RIFP-699 (1987);
H. Minakata, H. Nunokawa, S. Shiraishi and H. Suzuki, Contributed paper Nr. 85 and Tokyo University Preprint TMUP-HEL-8703 (1987).
- 65) M. Voloshin and M. Visotski, Preprint ITEP-1 (1986);
L. Okun, Preprint ITEP-14 (1986);
L. Okun, M. Voloshin and M. Visotski, Preprint ITEP-20 (1986).
- 66) F. Eisele, Rapporteur talk at this Conference.
- 67) R. Barbieri and G. Fiorentini, University of Pisa Preprint IFUP-TH-87/8 (1987).
- 68) M.J. Duncan, J.A. Grifols, A. Mendez and S.U. Sankar, Phys. Lett. 191 (1987) 304;
J. Liu, Phys. Rev. D35 (1987) 3447.
- 69) M. Fukugita and T. Yanagida, Contributed paper Nr. 195 and Tokyo University Preprint TU/87/307 (1987).
- 70) R.L. Gilliland, J. Faulkner, W.H. Press and D.N. Spergel, Ap. J. 306 (1986) 703.
- 71) S. Raby and G. West, Los Alamos preprints LA-UR-86-4152 (1986),
LA-UR-87-1255 (1987).
- 72) G.G. Ross and G. Segrè, Oxford University preprint 49/87 (1987).
- 73) G.B. Gelmini, L.J. Hall and M.J. Lin, Harvard preprint HUTP-86/A042 (1986).

- 74) H. Aihara et al., Contributed paper Nr. 346 (1987).
M.L. Perl, SLAC-PUB-4092 (1986) and to be published;
D.P. Stoker and M.L. Perl, JHU-HEP-870520, SLAC-PUB-4336 (1987).
- 75) E. Fernandez et al., Phys. Rev. D35 (1987) 1.
- 76) S. Raby and J. Valle, Private communication.
- 77) P. Avery et al., Contributed paper Nr. 258 (1987);
M. Artuso et al., Contributed paper Nr. 389 (1987).
- 78) E. Fishbach and S.A. Aronson, Proceedings of the XXIII International Conference on High Energy Physics, World Scientific, ed. S. Loker (1986) 1021, and references therein.
See also Ref. 36 and references therein.
- 79) P. Thieberger, Phys. Rev. Lett. 58 (1987) 1066.
- 80) C.W. Stubbs, E.G. Adelberger, F.J. Raab, J.H. Grundbach, B.R. Heckel, K.D. McMurry, H.E. Swanson and R. Watanabe, Phys. Rev. Lett. 58 (1987) 1070.
- 81) E.G. Adelberger et al., University of Washington preprint (1987).
- 82) I. Bars and M. Visser, Proceedings of the XIII International Conference on High Energy Physics, World Scientific, ed. S. Loker (1986) 1032.
- 83) R.D. Peccei, J. Solà and C. Wetterich, DESY preprint (1987).
- 84) A. de Rujula, CERN preprint TH. 4466 (1987).
- 85) J. Ellis, N.C. Tsamis and M. Voloshin, CERN preprint TH. 4719 (1987).