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by

R.D. Peccei

Deutsches Elektronen-Synchrotron DESY, Hamburg

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Looking Beyond the Standard Model

R.D. Peccei

Deutsches Elektronen-Synchrotron DESY,
Hamburg, Fed. Rep. Germany

ABSTRACT

I discuss a variety of theoretical, and some experimental, reasons for contemplating physics beyond the standard model. I illustrate this general discussion with two concrete examples. One concerns the GSI positron signals and their possible relation to variant axions. The other is connected with superstring theories (which I motivate in a heuristic manner) and the issue of neutrino masses in these theories.

I. Theoretical and Experimental Reasons to go Beyond the Standard Model

It is rather easy to make a long list of theoretical reasons for considering physics beyond the standard SU(3) x SU(2) x U(1) model of the strong and electroweak interactions. It is, however, very hard to find any experimental reasons to do so. Phenomenologically the standard model works marvelously well! Before discussing the few experimental open questions that there are, let me thus concentrate on the theoretical open questions. These questions are essentially related to structural aspects of the standard model and, instead of making a long list of queries, it proves useful to categorize them into three broad classes.

The first class of questions one can ask about the standard model is why is it a gauge theory based on the group SU(3) x SU(2) x U(1)? There are connected questions to this query. For instance, why at

$q^2 \sim 100\ \mathrm{GeV}^2$ are the strong and electroweak couplings so disparate in strength: $\alpha_3 \simeq 0.15$, $\alpha_2 \simeq 0.03$, $\alpha_1 \simeq 0.01$? Or why does the standard model (and we!) live in D=4 dimensions? Or where does gravity fit in the picture? Broadly speaking, we would like to have some idea really of why gravity, electromagnetism, the weak and the strong force are the forces we see in nature.

The second class of questions about the standard model is related to the kind of matter there is. We would really like to know why do the quarks and leptons transform as they do under the standard model. Of particular importance is trying to understand why do quarks and leptons appear in chiral asymmetric representations under the electroweak group? Related questions concern the U(1) assignments of the known fermions, which give third integral charge to the quarks but integral charge to the leptons. Another set of questions, of the same kind, concerns the reasons why a repetitive family structure appears and the issues of whether the number of repetitions stops at three and of whether families contain or not a right handed neutrino.

The third broad class of questions related to the structure of the standard model is concerned with the dynamics of mass generation. The spectrum of masses in the standard model is extensive and peculiar. The photon and gluon are massless and so are, for all intents and purposes, the neutrinos. The charged fermions we know, including a presumed top quark, range in mass over almost five orders of magnitude, from one half of an MeV to near 50 GeV. The W and Z masses are close to 100 GeV but we have as yet no evidence for where does the Higgs boson mass lie. There are associated questions here, too. For example, what physics fixes the Cabibbo Kobayashi Maskawa mixing angles? All elementary excitations we know get masses only after SU(2) x U(1) breakdown. That is

$$m \sim \Lambda_F \, , \qquad \text{(I.1)}$$

where $\Lambda_F = (\sqrt{2}\, G_F)^{-1/2} \simeq 250$ GeV is the scale of this breakdown. What fixes Λ_F and why is Λ_F so disparate from the other two scales we know:

$$M_{\mathrm{Planck}} \simeq 10^{19}\ \mathrm{GeV} \gg \Lambda_F \gg \Lambda_{QCD} \simeq 100\ \mathrm{MeV} \qquad \text{(I.2)}$$

where M_{Planck} is the scale associated with Newton's constant: $G_N \sim \frac{1}{M_{\mathrm{Planck}}^2}$ and Λ_{QCD} is the dynamical scale of QCD. Eq. (I.2) is related to the famous naturalness question of the scalar sector of the standard model /1/.

In addition to these three broad classes of theoretical questions concerning the structure of the standard model there is a related, perhaps more technical, question: Why is the $\bar\theta$ angle so small? This question is associated with the, so called, strong CP problem. Because of the nature of the QCD vacuum /2/ and the existence of a chiral anomaly /3/ the Lagrangian of the standard model is augmented by a T, P and CP violating term

$$\mathcal{L}_{CP\ viol.} = \frac{\alpha_3}{8\pi}\left[\theta + \mathrm{Arg}\det M\right]\tilde{F}_a^{\mu\nu} F_{a\mu\nu} = \frac{\alpha_3}{8\pi}\,\bar\theta\,\tilde{F}_a^{\mu\nu} F_{a\mu\nu} \qquad \text{(I.3)}$$

Here θ is the QCD vacuum angle and M is the quark mass matrix. The term in (I.3) can give rise to a nonvanishing electric dipole moment for the neutron and from the existing bound on this quantity /4/ one infers a bound on $\bar\theta$ /5/

$$\bar\theta \lesssim 10^{-8} - 10^{-9} \qquad \text{(I.4)}$$

This is one of two examples I know where there is a cancellation between two, apparently, distinct sectors of the standard model. For $\bar\theta \simeq 0$ it is necessary that the QCD contribution, θ, should cancel against the electroweak contribution coming from Arg det M. The renormalizability of the electroweak theory is insured by a similar amazing cancellation. Due to the U(1) assignment of quarks and leptons one has that

$$\mathrm{Tr}\, Q_{\mathrm{quarks}} + \mathrm{Tr}\, Q_{\mathrm{leptons}} = 0 \qquad \text{(I.5)}$$

thereby assuring that no chiral anomalies /3/ spoil the renormalizability of the SU(2) x U(1) theory. It would be nice to know a reason why these interrelations obtain in the standard model.

In comparison to the theoretical reasons for wanting to go beyond the standard model - the issues of forces, matter and mass dynamics - there is a dearth of experimental reasons to do so. I could think only of one good reason, and two more speculative reasons, for contemplating physics beyond SU(3) x SU(2) x U(1) (plus gravity). The good reason is related to the matter-antimatter asymmetry of the universe. We know that the ratio of baryon minus antibaryons in the Universe is approximately /6/

$$\frac{n_B - n_{\bar{B}}}{n_\gamma} \simeq 10^{-10} \ . \qquad \text{(I.6)}$$

Unless the Universe started that way - a peculiar initial condition - there must have been baryon number violating interactions at some level. However, in the standard model, baryon number is conserved (I will not worry here about possible non perturbative violation of baryon number in the standard model /2/ /7/). Hence to explain Eq. (I.6) it is necessary to have some physics beyond the standard model.

Less direct, but still quite compelling, evidence for physics beyond the standard model is related to neutrino masses and the solar neutrino problem, which has been the focus of much of this meeting. The discrepancy between the predicted rate /8/ for solar neutrinos in the Davis experiment /9/ and the measured rate:

$$\begin{aligned} R_{\mathrm{Theory}} &= 7.6 \pm 1.5 \ \mathrm{SNU} \\ R_{\mathrm{Measur}} &= 2.1 \pm 0.3 \ \mathrm{SNU} \end{aligned} \qquad \text{(I.7)}$$

could well be due to our imperfect understanding of the sun. It is very appealing, however, to suppose that this discrepancy is due to $\nu_e \to \nu_x$ conversion in the sun, as suggested recently by Mikheyev and Smirnov /10/ based on earlier work of Wolfenstein /11/. As Bethe has emphasized /12/, the MSW effect can occur even for small mixing angles, provided $\delta m^2_{\nu_e - \nu_x} \leq 5 \times 10^{-5}\ (\mathrm{eV})^2$, where ν_x is a distinct neutrino from ν_e. Barring accidental cancellations this would require a tiny neutrino mass

$$m_{\nu_x} \lesssim 7 \times 10^{-3}\ \mathrm{eV} \qquad \text{(I.8)}$$

Although a Dirac mass of this magnitude is not illegal in the standard model, it is silly! Much more reasonable is to suppose that such a small value for m_{ν_x} arises through the see-saw mechanism /13/:

$$m_{\nu_x} \sim \frac{m^2_{\ell x}}{M} \qquad \text{(I.9)}$$

where M is a large, lepton under violating, Majorana mass. If $x = \mu$ then Eq. (I.9) (which is a heuristic equation at any rate) implies $M \sim 10^9$ GeV, while for $x = \tau$ one has $M \sim 10^{12}$ GeV. Lepton number violating scales of these orders of magnitude obviously necessitate physics beyond the standard model.

The last potential evidence for physics beyond the standard model, which I would like to mention here, has been discussed by Schweppe /14/ at this meeting. This is related to the anomalous positrons, and correlated electrons, produced at GSI. This data can be understood, at the kinematical level, by supposing that what results in the heavy ion collision is a particle a, of mass near 1.7 - 1.8 MeV, which is produced nearly at rest and then decays into e^+e^- pairs. It should be mentioned that it is very difficult to understand /15/ how any such particle could be produced almost at rest in these collisions. Nevertheless, the fact that M_a is so light suggests that a might be an axion. Axions /16/ result from trying to solve the strong CP problem by imposing an extra symmetry /17/ on the standard model, which guarantees that $\bar{\theta} = 0$. Thus if the GSI phenomena has anything to do with axions there is the exciting possibility that we have found some physics beyond the standard model which actually explains one of the open puzzles. This would be progress. Alas, as I shall now discuss this does not turn out to be the case!

II. Variant Axions: Birth and Untimely Death

The strong CP problem can be solved by augmenting the standard model by a global chiral symmetry: $U(1)_{PQ}$. One can show /17/ that such a symmetry locks the phase of the quark mass matrix to the QCD vacuum angle θ , thereby leading to $\bar{\theta} = 0$. To impose a $U(1)_{PQ}$ symmetry on the standard model it is necessary to have two Higgs doublets, so that chiral quark rotations can be compensated by Higgs phase rotations. The Yukawa Lagrangian in the Peccei Quinn model /17/:

$$\mathcal{L}^{PQ}_{\text{Yukawa}} = \Gamma^{u}_{ij}\, \bar{Q}_{Li}\, \Phi_1\, u_{Rj} + \Gamma^{d}_{ij}\, \bar{Q}_{Li}\, \Phi_2\, d_{Rj} + h.c. \qquad \text{(II.1)}$$

where Q_{Li} are quark doublets of different families i, is clearly invariant under the chiral global $U(1)_{PQ}$ transformation:

$$\begin{aligned} u_{Rj} &\to e^{i\alpha} u_{Rj} & \qquad \Phi_1 &\to e^{-i\alpha}\Phi_1 \\ d_{Rj} &\to e^{i\beta} d_{Rj} & \qquad \Phi_2 &\to e^{-i\beta}\Phi_2 \end{aligned} \qquad \text{(II.2)}$$

(To rotate $\bar{\theta}$ away it is necessary that there be a net chiral transformation in the quark sector, so that $\alpha \neq -\beta$. /17/ It is for this reason that one needs two Higgs doublets, since with one Higgs doublet $\Phi_2 = \tilde{\Phi}_1 = i\tau_2 \Phi_1^{*}$ and thus $\alpha = -\beta$). Because of this extra symmetry, however, when the Higgs fields acquire vacuum expectation value:

$$\langle \Phi_i \rangle = \frac{1}{\sqrt{2}}\, f_i \qquad \text{(II.3)}$$

with

$$f = \sqrt{f_1^{\,2} + f_2^{\,2}} = (\sqrt{2}\, G_F)^{-1/2} = \Lambda_F \simeq 250 \text{ GeV} \qquad \text{(II.4a)}$$

being fixed by the scale of the SU(2) x U(1) breakdown and

$$x = f_2/f_1 \qquad \text{(II.4b)}$$

remaining a free parameter, also $U(1)_{PQ}$ breaks down. Associated with the breakdown of this global symmetry there arises a Goldstone boson - the axion /16/. In fact, since $U(1)_{PQ}$ is a chiral symmetry which suffers from an ABJ anomaly /3/ the axion is not really massless, but acquires a light mass /18/:

$$M_a \simeq m_\pi \frac{f_\pi}{f} \frac{(m_u m_d)^{1/2}}{(m_u + m_d)} N\left(x + \frac{1}{x}\right) \simeq 25\, N\left(x + \frac{1}{x}\right) \text{KeV} \qquad \text{(II.5)}$$

where N is the number of generations of fermions. Other properties of, what has been called, the standard axion are also fixed by the value of x. For instance the axion couplings to quarks is given by

$$\mathcal{L}_{aqq} = \frac{m_q}{f}\, \bar{q}\, i\gamma_5 q\, a \begin{Bmatrix} x \\ x^{-1} \end{Bmatrix} \qquad \text{(II.6)}$$

where the top line applies to charge 2/3 quarks and the bottom line to charge - 1/3 quarks.

The GSI excitation cannot be the standard axion since the standard axion has already been ruled out experimentally /19/! Actually, this analysis applies for an axion of mass less than $2m_e$ which has a very slow lifetime into two photons. It can, however, be repeated even in the case of an axion which decays rather rapidly into e^+e^- pairs. In particular, one of the processes $\psi \to \gamma a$ or $\Upsilon \to \gamma a$ badly violates the existing experimental bounds. One has /19/

$$B(\psi \rightarrow \gamma a) = (4.9 \pm 0.8) \times 10^{-5}\, x^{2} < 1.4 \times 10^{-5}$$
$$B(\Upsilon \rightarrow \gamma a) = (2.7 \pm 0.7) \times 10^{-4}\, x^{-2} < 3 \times 10^{-4} \qquad \text{(II.7)}$$

If $M_a \simeq 1.7 - 1.8$ MeV Eq. (II.5), for $N = 3$, tells us that x or x^{-1} is of order 20 and so one of the above decays would have given a spectacular signal, which has not been seen.

Although the quarkonia bounds are deadly for the standard axion, it is possible to consider a variant of the original Peccei Quinn model which neatly avoids these bounds. These variant axion models were originally discussed by Bardeen and Tye /18/ and were revived by Krauss and Wilczek /20/ and by Wu, Yanagida and myself /21/ as possible candidates for an explanation of the GSI phenomena. The idea behind these models is very simple. Instead of coupling Φ_1 to all u_{R_i} and Φ_2 to all d_{R_i} as in Eq. (II.1), in variant axion models one couples $\tilde{\Phi}_1$ only to some u_{R_i}, with the rest of the u_{R_i} being coupled to $\tilde{\Phi}_2$. Effectively these models have an $U(1)_{PQ}$ symmetry only for some doublets but not for all N of them. The simplest model /20//21/, for instance has only u_R coupled to $\tilde{\Phi}_1$, so that $N_{eff} = 1$, and only the u quark has a coupling proportional to x (cf. Eq. II.6). In view of Eq. (II.5), for $N_{eff} = 1$, x is very large ($x \simeq 70$). Furthermore, one can arrange the models so that also electrons couple to Φ_1 ($\tilde{\Phi}_1$ to be precise) and have an enhanced coupling leading to a fast decay lifetime:

$$\tau(a \rightarrow e^+e^-) = 3 \times 10^{-9}\, x^{-2} \text{ sec} \simeq 6 \times 10^{-13} \text{ sec} \qquad \text{(II.8)}$$

where the numerical value holds for a 1.7 MeV variant axion of the type described.

Variant axions appeared viable at the time of their proposal since they could escape the quarkonia bounds by construction and could, in view of their fast decay rate, avoid other bounds arising from beam dumps and nuclear deexcitation processes /19/. Because of the GSI signals they naturally elicited considerable interest. However, unfortunately,

these axions have also been ruled out experimentally now. Evidence against them, arising from an electron beam dump experiment at KEK, has been reported at this conference by Sasao /22/. I want to describe similar negative results arising from hadronic experiments, which really serve to rule out entirely any kind of variant axion model.

One can characterize variant axion models by their isovector and isoscalar mixing parameters λ_3 and λ_s, respectively /23/. (Axions do not carry in general definite isospin and the parameters λ_3 and λ_s detail their isospin couplings to hadrons). For the simplest variant model of Ref. /20/ and /21/ one has

$$\lambda_3 \simeq 3/8\, x \simeq 25 \; ; \quad \lambda_s \qquad \text{(II.9)}$$

One can show, however, that for <u>any kind</u> of variant axion model the difference between λ_3 and λ_s is fixed by the axion mass /23/:

$$(\lambda_3 - \lambda_s)^2 \simeq \left(\frac{3}{8} N x\right)^2 \simeq (25)^2 \qquad \text{(II.10)}$$

where the numerical value applies for a 1.7 MeV variant axion. There exist now data relevant for isovector axions which is in strong contradiction with the expectations of Eq. (II.9). Furthermore, there is also a bound on the isoscalar mixing parameter λ_s, which taken in conjunction with the bound on λ_3, is at variance with the result (II.10). Thus a combination of these bounds serves to rule out entirely all possible variant axion models.

The isovector bounds come from an experiment at Caltech /24/ which looked for axions in the deexcitation of the $2^+,1$ 9.17 MeV state of ^{14}N to its $1^+,0$ groundstate and from a search at SIN /25/ for axions in π^+ decay: $\pi^+ \rightarrow a\, e^+ \nu_e$. In both experiments axions of mass near 1.7 MeV decaying rapidly into e^+e^- pairs were looked for with negative results. Bardeen, Yanagida and I. /23/ have calculated the expectation of variant axion models for these decays and used the experimental results to bound λ_3. The axion to photon rate for the ^{14}N transition, using the results of Donnelly et al. /26/ to estimate the nuclear matrix element ratio, is

$$\frac{\Gamma_a}{\Gamma_\gamma} = 2 \times 10^{-5}(\lambda_3)^2 \qquad \text{(II.10a)}$$

to be compared to the experimental limit /25/

$$\frac{\Gamma_a}{\Gamma_\gamma} < 4 \times 10^{-4} \qquad \text{(II.10b)}$$

For the π^+ decay an even stronger bound on λ_3 is obtained. Theoretically /23/ one predicts

$$B(\pi^+ \to a e^+ \nu_e) = 3 \times 10^{-9}(\lambda_3)^2 \qquad \text{(II.11a)}$$

while the SINDRUM experiment gives

$$B(\pi^+ \to a e^+ \nu_e) < (1\text{-}2) \times 10^{-10} \qquad \text{(II.11b)}$$

(The range above depends slightly on the exact lifetime of the axion). Using Eq. (II.11) one has

$$\lambda_3 \lesssim 0.2 \qquad \text{(II.12)}$$

in strong disagreement with the prediction of the simplest variant axion model given in Eq. (II.9).

The bound on the isoscalar mixing parameter λ_s comes from a reanalysis by Calaprice et al. /27/ of an old internal pair conversion experiment of Warburton et al. /28/. The presence of axions decaying into e^+e^- pairs would have affected the measured angular distributions. For the isoscalar transition in ^{10}B from the 3.59 MeV $2^+,0$ state to the $1^+,0$ state at 0.72 MeV the reanalysis of Calaprice et al. gives the bound /27/

$$\frac{\Gamma_a}{\Gamma_\gamma} < 7.5 \times 10^{-5} \qquad \text{(II.13a)}$$

while our calculations yield

$$\frac{\Gamma_a}{\Gamma_\gamma} = 6.1 \times 10^{-4}(\lambda_s)^2 \qquad \text{(II.13b)}$$

Hence

$$\lambda_s \lesssim 0.35 \qquad \text{(II.14)}$$

Obviously, given Eq. (II.12) and Eq. (II.14), the constraint of Eq. (II.10) characteristic of variant axion models is badly violated and these models are ruled out. The possible remaining worry, connected with using an old experiment to bound λ_s, has actually now being removed by a direct bound /29/ on the ^{10}B transition of the 3.59 MeV $2^+,0$ state to its $3^+,0$ ground state. For this decay one predicts

$$\frac{\Gamma_a}{\Gamma_\gamma} = 7.9 \times 10^{-4}(\lambda_s)^2 \qquad \text{(II.15a)}$$

and the present bound at 90 % confidence level is /29/

$$\frac{\Gamma_a}{\Gamma_\gamma} < 7.2 \times 10^{-3} \qquad \text{(II.15b)}$$

giving

$$\lambda_s \lesssim 3.02 \qquad \text{(II.16)}$$

which, although not as strong as the bound of Eq. (II.14), suffices to rule out variant axions. My conclusion therefore is that the GSI signal remains mysterious but it has nothing to do with axions (or particle physics).

III. The Express Line to Superstrings

I want to return now to the central theoretical issues suggesting physics beyond the standard model - the questions of forces, matter and mass dynamics - and indicate, by means of a set of "non stringy" arguments, how the recently very popular superstrings theories /30/ give an extremely appealing answer to these questions. (In a subsequent section, though, I shall discuss also some of the troublesome issues in these theories.) The set of arguments I would like to put forward here involve three principal ideas. That of unification, that of supersymmetry and that of spontaneous compactification. Let me discuss these in turn.

The idea of unification of the existing forces into theories based on a simple gauge group /31/ is a natural extension of the process that led to the SU(2) x U(1) theory of electromagnetic and weak interactions. These, so called, GUT theories suppose that the standard model group is embedded in a larger group structure which, at some scale M_x, suffers a spontaneous breakdown, totally analogous to the breakdown of $SU(2) \times U(1) \rightarrow U(1)_{em}$ at Λ_F. The natural sequence for GUTs is the exceptional group sequence /32/, starting from the standard model $SU(3) \times SU(2) \times U(1) \equiv E_3$ to E_8, the largest exceptional group. This sequence is detailed in Table III.1.

Table III.1: The Exceptional Group Sequence

Group	Number of Gauge Fields
$E_3 \equiv SU(3) \times SU(2) \times U(1)$	12
$E_4 \equiv SU(5)$	24
$E_5 \equiv SO(10)$	45
E_6	78
E_7	133
E_8	248

With GUTs the question of forces changes from why SU(3)xSU(2)xU(1) to, say, why SU(5)? However, some important conceptual advantages have been gained. For instance, now at energies above M_x there is just one coupling constant α_u. The disparate values of α_3, α_2 and α_1 at low q^2 values obtain from the different renormalization group behaviour of the coupling constants below M_x /33/. Indeed, requiring that these

couplings unify at M_x into a single coupling α_u - a non trivial requirement - yields $M_x \sim 10^{14} - 10^{15}$ GeV. Hence the physics of GUT theories is characterized by energy scales much closer to the Planck scale.

If the forces are unified in a GUT group, there is also some simplification in the matter assignments, since it is necessary that the fermions of the theory themselves be in some GUT representation. These representations are much simpler than those under which the fermions transform in the standard model. Furthermore, since $G \supset SU(3) \times SU(2) \times U(1)$ one understands naturally why the particular U(1) assignments of the quarks and leptons (their charges, if you will) obtain. For instance, for SU(5) the fermions of each family sit in a $\bar{5}$ plus a 10 representation. For SO(10) all fermions of a given family are put into a 16, which requires the addition of a right handed neutrino $\nu_R \sim \nu_L^c$ to the theory for each family. In E_6 quarks and leptons sit in the 27, which requires that there be another 11 extra states, etc.

Although unification of the forces brings some simplification in the matter assignments, it brings no direct understanding of why we have in nature these particular matter representations. From the point of view of GUTs any fermion representation is ok for matter fields. To tie down the matter representations we need in some way to unify matter with forces. For that one needs supersymmetry, a symmetry which relates bosons to fermions /34/. To follow the line of logic I am adopting here, one wants the matter fermions to be partners of the gauge fields and so to sit in a vector multiplet V = (1, 1/2) of a supersymmetric GUT theory. This is not the assumption one makes normally in these theories. Usually matter sits in chiral multiplets $\Phi = (1/2, o)$ and gauge fields have their own separate gauginos. The pairing I am suggesting here in fact appears crazy as it runs into two fundamental problems:

i) Because the adjoint representation is real, pairing fermions to gauge fields in a vector multiplet V necessarily puts the fermions in real representations. However, we know that fermions are in chiral asymmetric representations!

ii) To get the quantum numbers of the known fermions right the GUT group must be at least E_6 and we need E_8 for 3 families. However, in all these cases there are many other extra fermions, whose traces are yet to be seen in nature!

Both of these points are illustrated by decomposing the 78 of E_6 into its SO(10) subgroup and the 248 of E_8 into its $E_6 \times SU(3)$ maximal subgroup. One has

$$78 = 45 + 1 + \overline{16} + 16 \qquad (III.1)$$

$$248 = (78,1) + (1,8) + (\overline{27},\bar{3}) + (27,3) \qquad (III.2)$$

The matter fields of one family can be put in the 16 of (III.1) and those of three families into the (27,3) of (III.2), including some additional states. However, one sees from the above that in addition to a 16 there is a $\overline{16}$ in (III.1) and analogously a $(\overline{27},\bar{3})$ in (III.2). So the fermions in these theories have chirally symmetric partners - mirror fermions - whose couplings with respect to the weak interactions are V+A. In addition both (III.1) and (III.2) contain many more real states besides the wanted matter fermions. The idea that force and matter are unified via supersymmetry makes no sense unless the above problems are somehow solved. Indeed, at this stage, also gravity is missing.

The third concept, that of spontaneous compactification, makes the above ideas potentially viable. In space time dimensions $D = 4n + 2$ it is possible to have fermions which are chiral and yet obey reality conditions /30/. Furthermore if the extra dimensions beyond four compactify, one may be able to obtain chiral fermions in four dimensions. The number of the produced chiral fermions in D=4 is related to the geometrical characteristics of the compactified space K. Essentially this number is the number of zero modes of the Dirac operator acting on the extra

dimensions /35/. If the properties of K are linked to the gauge fields present in the theory, and the theory is supersymmetric, then the chiral fermions that emerge in D=4 will have specific quantum numbers. Finally, since one is dealing with space time properties gravity is naturally included in these kind of considerations.

The above scenario is precisely what is thought to happen in superstring theories /30/. These theories are supersymmetric string theories in D=10 dimensions with a fixed gauge group which is either SO(32) or $E_8 \times E_8$ (Only this latter group is useful phenomenologically). It is not yet proven that superstring theories compactify to D=4, but if one assumes that this happens leaving an N=1 supersymmetry in four dimensions the compact manifold K is specified. This manifold, as Candelas, Horowitz, Strominger and Witten showed in a seminal paper /36/, is a so called Calabi-Yau space, which is Ricci flat, has a Kähler structure and has SU(3) holonomy. In particular, in these manifolds the gauge fields in an SU(3) subgroup of $E_8 \times E_8$ are identified with the spin connection, so that they are non vanishing. This means that in four dimensions the remaining gauge group is smaller than $E_8 \times E_8$. Assuming that this identification is made with respect to the SU(3) of the maximal $E_6 \times SU(3)$ subgroup of one of the E_8, the gauge structure present in 4 dimension is $E_6 \times E_8$. However, if K is not simply connected the E_6 group can be further broken by non vanishing flux loops /37/ to some subgroup of E_6. Since E_6 is one of the GUT groups one sees that superstrings offer the possibility of understanding why only certain forces appear. Note that in this respect the other E_8 provides a shadow world coupled to our world only gravitationally.

Superstrings offer also the possibility of understanding the matter content we see. Of the fermions present in the 248 adjoint representation of the non shadow E_8, only those that have a non trivial SU(3) content can emerge in four dimensions. Using Eq. (III.2) one sees that these are either in a 27 or a $\overline{27}$ of E_6. Now the 27 is precisely the right representation to contain the quarks and leptons! Furthermore, the number of 27 minus $\overline{27}$ which emerges is fixed by the geometry of K via an index theorem. One finds /36/ that

$$n_{27} - n_{\overline{27}} = 1/2\,|\chi(K)| \qquad \text{(III.3)}$$

where χ is the Euler number. Thus the number of families is determined by the topology of the Calabi-Yau manifold K, a remarkable result.

We see that even though in D=10 matter and forces were on equal footing, both sitting in the adjoint of $E_8 \times E_8$, in D=4 the emerging picture is rather disparate. Forces, acting on ordinary matter, are mediated by gauge fields which transform according to the 78 of E_6. Each of these gauge fields has an associated gaugino, so one has a vector multiplet V(1, 1/2) (There may be less gauge fields and gauginos if there is flux breaking). Matter on the other hand transforms as certain numbers of 27 and $\overline{27}$ of E_6, with the difference between 27 and $\overline{27}$ being fixed by the Euler number, Eq. (III.3). Each of these fermions will be accompanied by a complex scalar, since in the compactification the supersymmetry is preserved. Hence matter transforms as chiral multiplets $\Phi(1/2,0)$. Clearly not much symmetry is apparent in D=4 between forces and matter. However, in superstring theories both had the same origin.

There is another aspect of superstring theories which is also very nice and which I want to mention since it addresses the matter of mass dynamics discussed earlier. In principle all Yukawa couplings are fixed in the compactification. They originate basically from vertices involving fermions and gauge fields in D=10. Now gauge fields along the compact dimensions become scalar fields in D=4. Hence, very roughly speaking, the Yukawa couplings in D=4 are nothing but gauge couplings in D=10 after integrating over the compact space. Provided one can generate the Fermi scale Λ_F, superstring theories appear to be able to predict in principle all fermion masses, since these are just proportional to the Yukawa couplings times Λ_F. However, superstring theories really set up the physics at a scale M_{comp}, which should be of the order of M_{Planck}. To understand whether these considerations have anything to

do with reality one must try to make connection with physics at 100 GeV and see if there are any problems, or indeed if any of the predictions survive the transit.

IV. The Troublesome Neutrinos in Superstring Theories

I hope that my discussion in the last section has given you an idea of why many theorists find superstrings so appealing. Here is a theory which may provide the first example of a consistent quantum theory of gravity and which, at the same time, has within it all the elements to illuminate the structural mysteries of the standard model. This said, however, there remain profound difficulties to overcome in superstring theory before one can make a claim that it is at all connected to particle physics, as we know it. (There are also very hard problems to solve connected with the string aspects of superstrings, but these lie beyond both the scope of this talk and my own competence). Let me list a few of these difficulties:

i) With Calabi-Yau compactification the theory is left with an unbroken supersymmetry. Obviously this supersymmetry must somehow be broken but at scales close to Λ_F, if one wants to appeal to supersymmetry as a cure for the hierarchy problem /1/. How this breaking occurs and how Λ_F itself is generated is, however, not too clear. One appeals to the hidden sector of the theory to trigger the supersymmetry breaking /38/ and then to radiative corrections to generate Λ_F /39/, but no convincing demonstration, ab initio, of these mechanisms starting from the superstring theory has been given /40/.

ii) With flux loop breaking /37/ there are few workable models which survive at low energy. If there are no intermediate scales, one can show that these models must have an extra U(1) gauge group /41/. However, with intermediate scale breaking, it is possible that no such U(1) appears /41/. Hence it is difficult to pin down precise predictions arising from the superstring. Much of what can be seen at low energy seems to depend more on the path taken than on the original 10 dimensional theory.

iii) Superstring models have generic problems with both proton decay and with neutrinos, which require careful monitoring.

Let me discuss in some detail the last point above and state more clearly both the nature of the problem and some of the proposed solutions. I indicated earlier that superstring theories give definite predictions for Yukawa interactions (and also Higgs self couplings) in D=4. Although one cannot yet compute these terms in detail, the overall structure of these interactions is known. The superpotential (from which one can derive both the Yukawa and Higgs self interactions) contains trilinear products of 27's or $\bar{27}$'s, which have an E_6 invariant form, although their coefficients need not respect E_6 /41/:

$$W \sim (27)^3 + (\bar{27})^3 \qquad \text{(IV.1)}$$

Inside each 27 there is a full family of quarks and leptons, including a right-handed neutrino. In addition, there is an extra charge -1/3 quark and its antiquark, a lepton doublet plus its antidoublet and one more neutrino-like excitation. Specifically, decomposing the 27 in terms of its SU(5) content one has

$$27 = \{\bar{5} + 10 + 1\} + \{5 + \bar{5} + 1\} \qquad \text{(IV.2)}$$

where the $\{5 + \bar{5} + 1\}$ are the extra states. The usual matter, written in terms of left-handed fields, reads

$$\bar{5} + 10 + 1 = (d^c, L) + (u^c, e^c, Q) + \nu^c \qquad \text{(IV.3a)}$$

while the new matter reads

$$5 + \bar{5} + 1 = (g, H^c) + (g^c, H) + S \qquad \text{(IV.3b)}$$

Here, L,Q,H^c and H are all SU(2) doublets and g ist the extra charge -1/3 quark.

The presence of the extra 5 and $\bar{5}$ states in the 27 is the source of both the baryon number and neutrino problem. One can check that $(27)^3$ contains the following five terms involving g or g^c and the usual quarks and leptons /42/

$$QQg \qquad u^c d^c g^c \qquad QLg^c \qquad u^c e^c g \qquad d^c \nu^c g \qquad \text{(IV.4)}$$

If all these terms have nonvanishing coefficients then one has very rapid baryon number violation mediated by the scalar partners of g. Typically one presumes that these objects have mass of $O(\Lambda_F)$ and so the trouble is real! Two solutions have been proposed for this problem. The first assumes that the compactified manifold K is so clever that the actual coefficient for the first two, or the last three, terms in (IV.4) vanishes, thereby restoring Baryon number as a global (accidental) symmetry of the model. To my knowledge, however, no manifolds with this property have been exhibited. A more appealing solution is that there is an intermediate scale of breaking, below $M_{comp} \sim M_{Planck}$. In this case it might be possible to give sufficient high mass to the g quarks to avoid direct contradiction with experiment /41/ /43/.

The neutrino mass problem is of a similar nature. There are five potential "neutrino" helicity states per family in each 27, namely:

$$27: \ \nu, \nu^c, N, N^c, S \qquad \text{(IV.5)}$$

where N and N^c are the neutral leptons in the doublets H and H^c. Because of this it is quite natural to expect a Dirac mass for the neutrinos. Indeed the superpotential (IV.1) gives the following two terms involving the fields in (IV.5)

$$\nu \nu^c N^c \qquad NN^c S \qquad \text{(IV.6)}$$

and one sees that a Dirac mass for the neutrino field ν ensues if N has a vacuum expectation value. A possible solution to this problem is again to appeal for the vanishing, for topological reasons, of the first term in (IV.6). But again no examples exist where this really happens. Note that it is not possible here to appeal directly to the see saw mechanism /13/ to get small neutrino masses. To provide for a large Majorana mass for the right handed neutrino one needs a Higgs multiplet transforming as the 351 of E_6, but this multiplet is not part of theory. However, as suggested by Deredinger, Ibañez and Nilles /44/ and by Nandi and Sarkar /45/ such a term may effectively be generated from non renormalizable terms in the superpotential. For instance, a term like

$$W_{NR} \sim \frac{(27)(27)\,\overline{27}\,\overline{27}}{M_{comp}} \qquad (IV.7)$$

contains an effective Majorana mass for ν^c of order of the expectation value of $\frac{\langle\overline{27}\rangle^2}{M_{comp}}$. If there is an intermediate mass scale, so that some component of the $\langle\overline{27}\rangle$ is large, this effective Majorana mass may be sufficient to initiate the see-saw. Note that for this mechanism to be feasible there must be some $\overline{27}$ in the theory. Indeed only in these circumstances it is possible to contemplate the existence of intermediate mass scales because with only 27's the potential can never have a minimum, if there are large vacuum expectation values /41/.

A second possibility to obtain reasonable neutrino masses is to marry the neutrino field ν^c with some other E_6 singlet neutral fermions χ, an excitation which actually might be present after compactification /46/. (An analogous mechanism using gauginos was discussed by Mohapatra /47/). Then additional neutrino masses can arise from terms in the superpotential of the structure

$$W \sim 27\ \overline{27}\ X \qquad (IV.8)$$

Intermediate scale breaking triggered by the $\overline{27}$ vacuum expectation value can give large Dirac mass terms for the $\nu^c X$ combination. In conjunction with the usual $\nu\nu^c$ Dirac mass term one is left with a heavy Dirac neutrino and a massless Majorana neutrino. This latter field can then obtain a small mass from an induced supersymmetry breaking Majorana mass for the X /48/. Clearly this is just one of possibly many scenarios for generating neutrino masses. Unfortunately the relationship of these "solutions" of the neutrino problem to the original superstring theory

becomes quite remote.

V. Concluding Remarks

Given the state of theoretical uncertainties of what might really be the physics beyond the standard model, and the lack of any real experimental limits at present, all that one can really do now is to wait for new experimental developments. Important clues on whether the superstring ideas I discussed are on the correct track should emerge in the next five years, both from experiments at the accelerators now being completed (Tevatron, SLC, LEP and HERA) and from non accelerator experiments. Of particular importance here are solar neutrino experiments, which may well give us positive evidence for the existence of an intermediate mass scale. Accelerator experiments, on the other hand could discover supersymmetric partners of ordinary matter which are of fundamental importance for superstring theories (at least at the psychological level - if for nothing else!). However, it is my belief that theorists should not all migrate to the Planck mass. There are still unclear aspects of the standard model and it might well be that in trying to understand these - like the issue of strong CP - one might find the clue for what physics really lies beyond the standard model. Although the smart money is betting on supersymmetry and superstrings, I am not quite ready to bet that the meeting in this series in 1992 will change its name to Sneutrino 92.

References

/1/ For discussions see for example E. Witten, Nucl. Phys. B188 (1981) 513 or J. Ellis in Proceedings of the 1985 Lepton Photon Symposium, Kyoto, Japan

/2/ G. 't Hooft, Phys. Rev. D14 (1976) 3432

/3/ S. Adler, Phys. Rev. 177 (1969) 2426; J.S. Bell and R. Jackiw, Nuovo Cimento 60A (1969) 49; W.A. Bardeen, Phys. Rev. 184 (1969) 1848

/4/ I.S. Altarev et al., Phys. Lett. 102B (1981) 13

/5/ V. Baluni, Phys. Rev. D19 (1979) 2227;
R. Crewther, P. di Vecchia, G. Veneziano and E. Witten, Phys. Lett. 89B (1979) 123

/6/ See for example, M. Turner in Proceedings of the 1983 NATO ASI "Quarks, Leptons and Beyond", Munich, Fed. Rep. Germany

/7/ V.A. Kuzmin, V. Rubakov and M. Shaposhnikov, Phys. Lett. 155B (1985) 36

/8/ J.N. Bachall et al., Ap. J. 292 (1985) 279

/9/ R. Davis Jr. in Proc. of the Neutrino Mass Miniconference, Telemark, Wisconsin (1980)

/10/ S.P. Mikheyev and A. Smirnov, Nuovo Cimento 9C (1986) 17

/11/ L. Wolfenstein, Phys. Rev. D16 (1978) 2369

/12/ H. Bethe, Phys. Rev. Lett. 56 (1986) 1305

/13/ T. Yanagida in Proceedings of the Workshop on Unified Theory and Baryon Number of the Universe, KEK, Japan 1979; M. Gell Mann, P. Ramond and R. Slansky in Supergravity (North Holland 1979)

/14/ J. Schweppe, these proceedings

/15/ See for example J. Reinhardt et al., Phys. Rev. 33C (1986) 194

/16/ S. Weinberg, Phys. Rev. Lett. 40 (1978) 223;
F. Wilczek, Phys. Rev. Lett. 40 (1978) 279

/17/ R.D. Peccei and H.R. Quinn, Phys. Rev. Lett. 38 (1972) 1440;
Phys. Rev. D16 (1977) 1791

/18/ W.A. Bardeen and S.-H.H. Tye, Phys. Lett. 74B (1978) 229

/19/ For a review, see for example A. Zehnder, Proceedings of the 1982 Gif-sur-Yvette Summer School

/20/ L.H. Krauss and F. Wilczek, Phys. Lett. 173B (1986) 189

/21/ R.D. Peccei, T.T. Wu and T. Yanagida, Phys. Lett. 172B (1986) 435

/22/ N. Sasao, these proceedings

/23/ W.A. Bardeen, R.D. Peccei and T. Yanagida, DESY 86-054, Nucl. Phys. B to be published

/24/ M.J. Savage et al., Phys. Rev. Lett. 57 (1986) 178

/25/ R. Eichler et al., SIN Preprint PR 86-07

/26/ T.W. Donnelly et al., Phys. Rev. D18 (1978) 1607

/27/ F.P. Calaprice et al., Internal Princeton Report 1986

/28/ E.K. Warburton et al., Phys. Rev. B133 (1964) 42

/29/ F.W.N. de Boer et al., Groningen preprint, Phys. Lett. B to be published

/30/ For reviews see J.H. Schwarz, Phys. Rep. 89C (1982) 233;
M.B. Green, Surveys in High Energy Physics 3 (1982) 127

/31/ For a review see G.G. Ross Grand Unified theories (Benjamin, San Francisco 1985)

/32/ R. Slansky, Phys. Rep. 79C (1981) 1

/33/ H. Georgi, H.R. Quinn and S. Weinberg, Phys. Rev. Lett. 33 (1974) 451

/34/ For a review see J. Bagger and J. Wess, Supersymmetry and Supergravity (Princeton Univ. Press, 1984)

/35/ For a discussion, see for example, P. Candelas, G.T. Horowitz, A. Strominger and E. Witten, in Proceedings of the Argonne Symposium on Anomalies, Geometry, and Topology (World Scientific 1985)

/36/ P. Candelas, G. Horowitz, A. Strominger and E. Witten, Nucl. Phys. B258 (1985) 46

/37/ Y. Hosotani, Phys. Lett. 126B (1983) 309; 129B (1983) 193;
E. Witten, Phys. Lett. 149B (1984) 351

/38/ M. Dine, R. Rohm, N. Seiberg and E. Witten, Phys. Lett. 156B (1985) 55; J.P. Derendinger, L.E. Ibañez and H.P. Nilles, Phys. Lett. 155B (1985) 65

/39/ L.E. Ibañez, Phys. Lett. 118B (1982) 73; K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, Prog. Theor. Phys. 68 (1982) 927

/40/ M. Dine and M. Seiberg, Phys. Lett. 162B (1985) 299

/41/ E. Witten, Nucl. Phys. B258 (1985) 75; M. Dine, V. Kaplunovsky, M. Mangano, C. Nappi and N. Seiberg, Nucl. Phys. B259 (1985) 549

/42/ J.D. Breit, B. Ovrut and G. Segré, Phys. Lett. 158B (1985) 33

/43/ F. del Aguila, G. Blair, M. Daniel and G.G. Ross, CERN TH 4336 (1985), Nucl. Phys. to be published

/44/ J.P. Derendinger, L.E. Ibañez and H.P. Nilles, CERN TH 4228 (1985), Nucl. Phys. to be published

/45/ S. Nandi and U. Sarkar, Phys. Rev. Lett. 56 (1986) 566

/46/ E. Witten, Princeton preprint (1985)

/47/ R.N. Mohapatra, Phys. Rev. Lett. 56 (1986) 561

/48/ R.N. Mohapatra and J. Valle, Maryland preprint 86-127