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THEORETICAL SURVEY OF HIGGS BOSONS AND AXIONS*

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ABSTRACT

The success as well as the problems of the minimal Standard Model are recalled. We survey essentially this Model and the theory of the standard axion (Nambu–Goldstone boson). Possible invisible and visualized (theoretical) axions are discussed as are certain astrophysical aspects of the existence of an axion. We survey also axion cosmology in superstring models and its consequences in the new anomaly cancellation mechanism to the sense of Green and Schwarz. Recent results for the search of the Higgs boson, and the axion are resumed. A great important is reserved for discussion of the Standard Model.

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1. INTRODUCTION

The Standard Model [Glashow (1961), Weinberg (1967), Salam (1968)] represents the most successful description developed to date of the physical world. None is so direct as the discovery of the boson W and then of the Z (1983) by the UA1 and UA2 Collaborations at CERN. Values for M_W and M_Z have been measured: $M_W = (80.8 \pm 2.7)$ GeV/c^2 and $M_Z = (99.8 \pm 1.6)$ GeV/c^2. The ratio M_W/M_Z calculated with these values agrees well with that given by $\cos\theta_W$. (The angle θ_E is usually expressed as $\sin^2\theta_W$ and is measured in neutrino-scattering experiments to be $\sin^2\theta_W = 0.224 \pm 0.015$). These results are not yet quite at the level of testing the electroweak radiative corrections at the one loop level.

The existence of nonstandard Higgs bosons (charged, very light or even massless) would open up some new opportunities for an experimental search for these particles. The evolution has been so decisive that at this time one could expect to perform a precision test of the Standard Model. New particles could be discovered, entirely new ideas related to new gauge theories, symmetry breaking, and supersymmetry might become relevant to the physics. It has pointed up [Jilles (1984), Haber and Kane (1985)] that in all supersymmetric models, physical charged bosons are essentially predicted.

We will be unable to cover such a vast subject. Instead of reviewing these successes, we will concentrate on a few possible problem areas: 1) The Higgs boson from a search for hypothetical particles in data from PEP [Feldman (1985)], TASSO]Komamiya (1986)] and CRYSTAL BALL [Lowe (1986)]. The CP–violation from possible astrophysical bounds on the axion mass in data from Cowan *et al.* (1986) and Tsai (1986) (laboratory experiments). The axionic anomaly (a new anomaly cancellation mechanism to the sense of Green and Schwarz (1984)] for axion cosmology in superstring models.

We will be unable to cover such a vast subject. We will survey essentially things that either were not covered in previous reviews or for which something new can be added.

2. The Minimal Standard Model

Apart from gravity, the minimal Standard Model successfully accounts for all known physics. Nevertheless, it cannot be regarded as the ultimate theory. However, some problems and shortcomings are present. We shall describe two of these problems only; that is: problems with the Higgs boson, and the strong CP–violation problem [Herczeg, Hoffman (1986)], together with some of the proposed theoretical schemes they motivate.

2.1 Problems with the Higgs Boson

The spontaneous breakdown of the electroweak gauge symmetry, necessary to generate the masses of the W, Z, and of the charged fermions, is implemented in the minimal Standard Model through the inclusion of elementary scaler fields. One objection to this approach is that it introduces a large number of parameters into the model. Another is a problem associated with the mass of the Higgs boson. The square of the zeroth order Higgs mass receives quadratically divergent contributions from radiative corrections. It is expected that the divergent term is cut off by the next largest mass scale in the model. For the minimal Standard Model, the only further mass scale is the Planck mass $m_P \simeq 10^{19}$ GeV ($m_P = (G_n)^{-1/2}$, G_N = Newton's gravitational constant). The result is that in order to obtain a

Higgs mass not larger than about 1 TeV (required if the weak interactions are not to become strong at energies above $\sim$ 1 TeV), a fantastic cancellation is needed between the zeroth order mass and the radiative corrections.

There are several proposed solutions to this problem [Langacker (1981), Peskin (1985)].

The minimal Standard Model requires the existence of only one neutral Higgs boson: The doublet $\phi = \begin{pmatrix} \psi^+ \\ \psi^0 \end{pmatrix}$ with hypercharge $y = 1$. (In nonminimal Standard Models, there could be more than one SU(2) doublet Higgs field or even SU(2) triplet Higgs fields. For example, in minimal supergravity models, two SU(2) doublet Higgs fields are expected [Haber and Kane (1984); Ellis *et a.*, (1985)]. In this case there would be two charged physical Higgs bosons ($H^{\pm}$) and three neutral ones (h_1^0, h_2^0 and h_3^0)).

Write the Higgs field as [Ansel'm, Ural'tsev and Khoze (1985)]:

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} = \begin{pmatrix} \phi^+ \\ \frac{1}{2}\left(v + x + 1 + ix_2\right) \end{pmatrix} \tag{1}$$

$$v \neq 0 \, , \qquad \langle x_1 \rangle = \langle x_2 \rangle = 0$$

In the case of a spontaneous development of a nonvanishing vacuum expectation value ($VEV = v \neq 0$) of the field $\phi(V = (1/\sqrt{2})\langle 0|\phi|0\rangle)$, the SU(2) $\otimes$ U(1) gauge group is broken down to the $\mathrm{U}(1)_{em}$ group. The W and Z bosons and the fermions acquire masses

$$M_W^2 = g^2 \, \frac{V^2}{4} \, , \quad M_Z^2 = g_2^2 \, \frac{V^2}{4} \, , \quad M_f = h_f \, \frac{V\sqrt{2}}{2} \, , \tag{2}$$

$g_2^2 = g^2 + g^{12}$ ($g = g_2 \cos\theta_W$, $g' = g_2 \sin\theta_W$), h_f are the Yukawa constants. The scalar fields ϕ^+ and x disappear from the spectrum of physical states because of the Higgs mechanism. Only the single neutral scalar boson H^0 turn out to be observable. Using Eq. (2) we can write the interaction of H with the W and Z bosons in the following form

$$\begin{aligned} & g^2\,\frac{\lambda}{2}\,W^+_\mu W^\mu H \;+\; \frac{g^2}{4} W^+_\mu W^\mu H^2 \\ & + \; g^2\,\frac{\lambda}{4\cos^2\theta_W}\;Z_\mu Z^\mu H + \frac{g^2}{8}\;\cos^2\theta_W X_\mu Z^\mu H^2 + \ldots \;. \end{aligned} \tag{3}$$

Under this form the Lagrangian in the Higgs sector reflects a basic property of Higgs bosons: their interaction with particles is proportional to the mass of these particles. In the fermion case, the mass appears in the amplitude, while in the boson case the square of the mass appears in the amplitude.

In a minimal Standard Model with a single Higgs doublet, the V can be expressed unambiguously in terms of the Fermi constants G_F with the help of Eq. (2):

$$V = \left(G_F\sqrt{2}\right)^{-1/2} \sim 146 \text{ GeV} \;. \tag{4}$$

The minimal neutral Higgs bosons can be mainly produced in the decay of heavy quarkonium [Wilckek (1978)] accompanied by a monochromatic photon.

The branching ratio of the upsilon decaying into a minimal Higgs plus a photon is given by [Komamiya(1985)]

$$B(\Upsilon \to \gamma H^0) = B(\Upsilon \to \mu\mu)\,\frac{G_F m_b^2}{\sqrt{2}\,\pi\alpha}\,\left(1 - \frac{m^2 H^0}{m^2 \Upsilon}\right) , \tag{5}$$

which is about $\mathcal{O}(10^{-4})$ if the Higgs boson mass is small.

The decay width of the minimal Higgs boson into a fermion pair $f\bar{f}$ is given by

$$\Gamma(H^0 \to f\bar{f}) = \frac{G_f m_f^2 m_{H^0}}{4\pi\sqrt{2}} \left(1 - \frac{4m_f^2}{m_{H^0}^2}\right)^{2/3} . \tag{6}$$

We consider now a Standard Model [Ansel'm, Ural'tsev and Khoze (1985)] with n doublets

$$\phi_i = \begin{pmatrix} \phi_i^+ \\ \frac{1}{\sqrt{2}\,(V_i + X_i^0 + iY_i^0)} \end{pmatrix} , \tag{7}$$

$$\langle X_i^0 \rangle = \langle Y_i^0 \rangle = 0 .$$

This model requires that no more than three doublets [Glashow and Weinberg (1977)] (ϕ_{11}, ϕ_2 and ϕ_3, say) initially interact with fermions. The most general form of their interaction would be

$$\mathcal{L}_{W+} = \sum_1^3 \pi(-1)^i \left(\frac{\phi_i^+}{V_i}\right) m_{\alpha_i} + h.a. , \tag{8}$$

for charged bosons, or

$$\mathcal{L}_{Z_0} = \sum_1^3 \pi \left(\frac{1}{V_i}\right) m_{\beta_i} \left(x_i^0, j_i^0\right) , \tag{9}$$

for neutral bosons, where

$$
\begin{aligned}
m_{\alpha_i} &= m_d\,\bar{u}'_L\,d_R + m_s\,\bar{c}'_L\,s_R + m_b\,\bar{t}'_L\,b_R \\
&\quad + m_u\,\bar{u}_R\,d'_L + m_c\,\bar{c}_R\,s'_L + m_t\,\bar{t}_R\,b'_L \\
&\quad - \left(m_e\,\bar{\nu}_e e_R + m_\mu\,\bar{\nu}_\mu\,\mu_R + m_t\,\bar{\nu}_t\,\tau_R\right) , \\
m_{\beta_i}(x_i^0, y_i^0) &= (m_d\,\bar{d}d + m_s\,\bar{s}s + \ldots)\,x_1^0\;(m_d\,\bar{d}\,i\,\gamma_5\,d + m_s\,\bar{s}\,i\,\gamma_5\,s + \ldots)\,y_1^0 \\
&\quad (m_u\,\bar{u}u + m_c\,\bar{c}c + \ldots)\,x_2^0\;(m_u\,\bar{u}\,i\,\gamma_5\,u + m_c\,\bar{c}\,i\,\gamma_5\,c + \ldots)\,y_2^0 \\
&\quad (m_e\,\bar{e}e + m_\mu\,\bar{\mu}\mu + \ldots)\,x_3^0\;(m_e\,\bar{e}\,i\,\gamma_5\,e + m_\mu\,\bar{\mu}\,i\,\gamma_5\,\mu + \ldots)\,y_3^0 \,.
\end{aligned}
$$

The vacuum expectation values V_i and the components ϕ_i^+ X_i^0 and y_i^0 are defined in Eq. (1). In Eq. (8) we have

$$
\begin{pmatrix} d' \\ s' \\ b' \end{pmatrix}_L = \begin{pmatrix} d \\ s \\ b \end{pmatrix}_L \quad , \quad \begin{pmatrix} u' \\ c' \\ t' \end{pmatrix}_L = \begin{pmatrix} u \\ c \\ t \end{pmatrix}_L \quad , \qquad (10)
$$

where the $U_i j$ are the elements of the standard K–M matrix [Kobayashi and Maskawa (1973)].

In order to use Eqs. (8) and (9) we need to know, in addition to the ratios of the VEV's V_i, how the physical Higgs bosons with a certain mass are constructed from the fields ϕ_i^+, x_i^0 and y_i^0.

The structure of interactions (8) and (9) is reproduced in an arbitrary Higgs sector under the sole requirement of natural conservation of flavors in the exchange of neutral scalar particles.

We should emphasize that both charged and neutral scalar particles also appear in technicolor models: various pseudo–Goldstone bosons. All interactions—both gauge and Yukawa interactions—of charged pseudo–Goldstone bosons are the same as the interactions of the elementary Higgs bosons, and in general they have only certain specific symmetry limitations.

Charged Higgs bosons can be found in e^+e^-annihilation since they are pair produced from a virtual photon. Since charged Higgs bosons normally couple to heavy fermions such as τ, c and b, the dominant decay modes are $H^- \rightarrow \tau^-\bar{\nu}_\tau$, $H^- \rightarrow s\bar{c}$, and $H^- \rightarrow b\bar{c}$. These dominant decays and also the H^- branching fractions have been studied by several experimental groups as they are shown in Table I.

In the last few years an extensive experimental search for supersymmetric particles has been carried out at PEP and PETRA (1985) to achieve a more fundamental understanding of the Higgs sector. The charged Higgs bosons which appear in the model described above are of interest from experimental standpoint [Ioganse n, Uralsev and Khoze (1982)].

3. THE AXION

3.1 The Strong CP–Violation Problem

Due to nonperturbative instanton effects, the Lagrangian of QCD contains a term that violates both parity and time-reversal invariance. This term is proportional to a parameter θ that is made up of two unrelated contributions: $\theta = \theta_{QCD} + \theta_{weak}$. The parameter θ_{QCD} resides in QCD, while θ_{weak} comes from the nonstrong section of the minimal Standard Model. The problem is how to make $\theta = \theta_{QCD} + \theta_{weak}$ as small as $10^{-8} - 10^{-4}$, an upper bound dictated by the experimental limit on the dipole momentum of the neutron ($d\mu < 3.6 \times 10^{-25}$ e-cm) [Altarev, Borisov, Borovikova *et al.* (1981)]. One solution is to suppose that the total Lagrangian (QCD + electroweak) possesses a global U(1) symmetry. The P,T–violating term can be then transformed to zero. The extra U(1) symmetry can be imposed in the Standard Model by extending the Higgs sector to contain at least two Higgs doublets. The spontaneous breakdown of the electroweak symmetry breaks the U(1) symmetry implying the existence of a Goldstone boson—the axion. The axion is not massless, but acquires a mass through an anomaly.

The simplest axion model (with two Higgs doublets and the U(1) symmetry broken at the weak scale) now appears to be ruled out experimentally, but more complicated axion models are possible.

We review here the standard axion and the visible and invisible axion, the basic implications between them, and their constraints from astrophysics and cosmology.

The Standard Axion. The standard axion, now known as the Peccei, Quinn, Weinberg and Wilczek axion, is a Nambu–Goldstone pseudoscalar boson with a mass of the order of a few hundred KeV ($> 200\ KeV$), spinless ($\mathcal{O}^-$) and a lifetime τ [Weinberg (1978), Bardeen and Tye (1978)], $\tau(a \to \gamma\gamma) \sim (0.8)(100\,\mathrm{KeV}/m_a)^5$ sec.

1. Position of the Problem.

Let us first consider the SU(2) $\otimes$ U(1) theory with two Higgs doublets. Peccei and Quinn (1977) noted that this theory does not have mass terms and hence there exists a possibility of introducing $\mathrm{U(1)_A}$ symmetry in GWS Lagrangian. It is possible because there exist complex Higgs fields.

For two Higgs doublets, ϕ_1, and ϕ_2, the most general Yukawa couplings to quarks and Higgs potential is [Kim (1982)]

$$
\begin{aligned}
\mathcal{L}_y = & - f^u_{ij}\left(\bar{q}_{Li}\,\phi_2\,u_{Rj} + \phi_2^+\,\bar{u}_{Ri}\,q_{Lj}\right) - f^d_{ij}\left(\bar{q}_{Li}\,\phi_1\,d_{Rj} + \phi_1^+\,\bar{d}_{Ri}\,q_{Lj}\right) \\
& - h^u_{ij}\left(\bar{q}_{Li}\,\phi_1^*\,u_{Rj} + \tilde{\phi}_1^+\,\bar{u}_{Ri}\,q_{Lj}\right) - h^d_{ij}\left(\bar{q}_{Li}\,\phi_2^*\,d_{Rj} + \tilde{\phi}_2^+\,\bar{d}_{Ri}\,q_{Lj}\right) \\
& - V\left(\phi_1,\phi_2\right) + \mu^2 h\left(\phi_1^+\phi_2^* + \tilde{\phi}_2\phi_1\right) \\
& - \sum_{i\neq j} b_{ij}\,\phi_i^+\,\phi_j^*\,\tilde{\phi}_j\,\phi_i - \sum_{i\neq j} c_{ij}\,\phi_i^+\,\phi_j^*\,\phi_i^+\,\phi_j^*\ ,
\end{aligned} \tag{11}
$$

where Higgs potential has the form [Grzadkowski (1986)].

$$
\begin{aligned}
V(\phi_1,\phi_2) = & \ \mu_1^2\phi_1^+\phi_1 + \mu_2^2\phi_2^+\phi_2 \\
& + \lambda_1\left(\phi_1^+\phi_1\right)^2 + \lambda_2\left(\phi_2^+\phi_2\right)^2 \\
& + \lambda_3\left(\phi_1^+\phi_1\right)\left(\phi_2^+\phi_2\right) + \lambda_4\left(\phi_1^+\phi_2\right)\left(\phi_2^+\phi_1\right) \\
& + \lambda_5\left[\left(\phi_1^+\phi_2\right)^2 + \left(\phi_2^+\phi_1\right)^2\right]\ ,
\end{aligned} \tag{12}
$$

after discrete transformation:

$$
\begin{array}{lcl}
\phi_1 \to \phi_1 & & \phi_1 \to -\phi_1 \\
 & \text{or} & \\
\phi_2 \to -\phi_2 & & \phi_2 \to \phi_2
\end{array}\ , \tag{13}
$$

which assures us of a natural flavor conservation at the tree level; $q)Li$ is the ith family quark doublet, b is a real symmetric and c is hermitian. The Higgs doublets are

$$\phi_1 = \begin{pmatrix} \phi_1^+ \\ \phi_1^0 \end{pmatrix} \quad , \quad \phi_2 = \begin{pmatrix} \phi_2^0 \\ \phi_2^- \end{pmatrix} \quad .$$

The Lagrangian, Eq. (11), do not have the desired $\mathrm{U(1)_A}$ symmetry. However, if $h_{ij}^u = h_{ij}^d = 0$, $\mu_h^2 = 0$ and $c_{ij} = 0$, a symmetry–the so-called Peccei–Quinn G–appears:

$$\begin{aligned} G: \quad u_{L,R} &\to e^{\pm i\frac{\alpha}{2}}\, u_{L_1R} \\ d_{L,R} &\to e^{\pm i\frac{\alpha}{2}}\, d_{L_1R} \\ \phi_{1,2} &\to e^{i\alpha}\, \phi_{1,2} \quad . \end{aligned} \tag{14}$$

Thus the Peccei–Quinn solution is natural. Therefore the imposed invariance G is

$$\mathrm{S} = \mathrm{SU(2)} \otimes \mathrm{U(1)} \otimes \mathrm{U(1)_A} \; . \tag{15}$$

2. Consequences.

By this symmetry mechanism,

(a) We can freely introduce the weak CP violation as far as the CP violation mechanism does not remove $\mathrm{U(1)_A}$ of Eq. (15).

(b) $\mathrm{U(1)_A}$ must be spontaneously broken since there does not exist a massless quark. Both $\langle\phi_1\rangle$ and $\langle\phi_2\rangle$ should develop nonvanishing vacuum expectation values. Since the continuous symmetry is spontaneously broken, there appears a Goldstone boson (an axion) [Wilczek (1978), Weinberg (1978)].

(c) This axion acquires a very small mass due to chiral symmetry breaking [Callan, Dashen and Gross (1976), Jackiw and Rebbi (1976), 't Hooft (1976)] by instantons.

In general , we obtain the axion field from the phase fields of $U(1)_G$ nontrivial Higgs fields

$$a = \frac{1}{N^{1/2}} \sum_i V_i G_i P_i$$
$$n = \sum_i V_i^2 G_i^2 , \tag{16}$$

where $P - i$ are the phase fields of Higgs fields ϕ_i. Depending on the value N, the resulting axion can be visible or invisible.

In general, a useful tool for distinguishing the axion visibility consists in using the following formula [Kim (1982)] of the axion coupling to matter fields

$$-ia \sum_i \frac{m_i}{N^{1/2}} \bar{f}_i \gamma_5 f_i (G_{iL} - G_{iR}) , \tag{17}$$

where, $m_i = f_{ij} V_i G(\phi_j^0)$, f_{ij} is the Yukawa coupling matrix of the fermion f_i and the scalar ϕ_j^0.

The Visible Axion

The visible axion is based on the standard axion from the Peccei–Quinn, Weinberg and Wilczek symmetry mechanism, and on the formula for distinguishing the axion visibility, Eq. (17). The visible axion occurs when (light quark mass)$/N^{1/2}$ is a Yukawa coupling strength. Thus the axion mass to the Bardeen–Tye (1978) is

$$m_a^{BT} = m_\pi \frac{f_\pi}{f_a} \left(x + \frac{1}{x}\right) N_g \frac{\sqrt{z}}{1+z} . \tag{18}$$

N_g is the generation number of fermions, $x = V_1/V - 2$, $z = m - u/m_d$, $m_\pi^2 f_\pi^2 = -m_u\langle\bar{u}u\rangle - m_d\langle\bar{d}d\rangle$. In fact, z is given by tree level current algebra to be $z \simeq 0.56$ [Weinberg (1977), Kaplan (1985)]. $V = f_a/\sqrt{2}$ [Sikivie (1982)], $f_\pi/f_a \ll 1$ for $F_\pi = 93$ MeV of QCD.

The Invisible Axion

The standard axion does not exist. But we can obtain the invisible axion when the Peccei–Quinn symmetry is broken at grand unification scales.

The invisible axion become visible axion by different mechanism of symmetry [Kim (1979), Schifman, Vainstein and Zahkavov (1980), Chikashige, Mohapatra and Peccei (1980), Dine *et al.*, Wise *et al.* (1981)]. Thus, for example Kim (1980), using Bardeen–Tye, and Dine, Fischler and Srednicki symmetry mechanisms, obtains a new mechanism to find the axion mass. The result is that

$$\begin{aligned} m_a^Q &= \frac{z^{1/2}}{1+z}\,\frac{f_\pi}{\widetilde{V}}\,\frac{m_\pi \alpha_C^2}{\pi^2}\,\ell n\,\frac{m^2 Q}{m_u m_d} \\ m_a^{DFS} &= \frac{\sqrt{2}}{\widetilde{V}}\,f + \pi\, m_\pi\, N_g\,\frac{z^{1/2}}{1+z}\,, \end{aligned} \tag{19}$$

where $\widetilde{V}$ is the scale where $\mathrm{U(1)_G}$ is broken ($\widetilde{V} \sim 10^7$ GeV for the heavy quark invisible axion (m_a^Q) if the heavy quark is the only contribution to the axion mass). We obtain for $m_a^{DFS} \sim 2.5 \times 10^{-8}$ MeV, for $\widetilde{V} = 10^9$ GeV.

There are other possibilities of realizing Goldstone bosons in Nature, satisfying the rules of Georgi *et al.* (1981), Kim (1981), Claudson *et al.* (1981), and Frampton and Kephart (1982).

The Role of the Axion in Cosmology and Astrophysics There are several astrophysical contributions for the Peccei–Quinn symmetry for obtaining the bound of the axion mass. Table II resumes these results and Fig. 2 the Feynman graphs for axion emission processes.

NOTE THAT:

(*a*) If the mass scale $\widetilde{V}$ is the electroweak scale -250 GeV, the resulting axion is visible.

(*b*) If the $U(1)_A$ breaking scale is much larger than this scale, by giving a vacuum expectation value to a $SU(2) \times U(1)$ singlet Higgs field, the resulting axion is invisible.

Since the astrophysical bounds on $\widetilde{V}$ is in the intermediate mass scale 10^7 GeV - 10^9 GeV, one may try to introduce an intermediate mass scale $\widetilde{V}$ in a grand unified theory so that the bad feature of the invisibility is made visible in the stars.

The cosmic string theory [Bennett (1982); Vilenkin, Turok, Kibble, Witten (1985); Bennett, Brandenberger and Truok, and Bagger (1986)] of galaxy formation has recently begun to attract some attention. The possible astrophysical role of superstrings is suggested by Witten (1985) by showing that the superstring (a superstring at rest and running in the x_3 direction. It is surrounded by a curve γ which bounds a surface S as shown in Fig. 0) is an axion string, the boundary of an axion domain wall (confinement of superstrings).

To show that this long string is an axion string, following Witten, we pick a contour γ that circles the string and calculate the changes in ϕ (the field of an axion) in circling the string,

$$\Delta = \int_{\gamma} d \sum^{\mu} \frac{\partial \phi}{\partial x^{\mu}} \,. \tag{20}$$

The string is an axion string if $\Delta \neq 0$.

One axion is model independent and arises as follows. Consider components of B_{MN} (antisymmetric tensor) with ordinary space-time M^4 indices $B_\mu \nu$, $\mu\,\nu = 0, 1, 2, 3$. The field strength corresponding to $B_{\mu\nu}$ is $H_{\mu\nu\alpha}$ which, in vacuum, satisfies the field equation $\partial^\mu H_{\mu\nu\alpha} = 0$. If we define the dual of $H_{\mu\nu\alpha}$ as $Y^\mu = (1/6)\,\epsilon^{\mu\nu\alpha\beta}\, H_{\nu\alpha\beta}$, we see that the field equation implies

$$\partial^\mu Y^\nu - \partial^\nu Y^\mu = 0 \Longrightarrow Y^\mu = \partial^\mu \theta$$

for some field θ. Δ depends on $\partial^\mu H_{\mu 30}$, *i.e.*, by Gauss's law and taking S to lie in the $x_1 x_2$ plane, we have

$$\Delta = \frac{1}{2} \int_S dx_1\, dx_2\, \partial^\mu H_{\mu 30} = \frac{1}{2}\, g \int_S dx_1\, dx_2\, V_{30} \ , \tag{21}$$

since $\partial^\mu H_{\mu\nu\alpha} = g\, V_{\nu\alpha}$, where g is a coupling constant, and V is the vertex operator. Thus we must show that V_{30} has a nonvanishing expectation value in the vicinity of the string.

This follows from the definitions of the vertex operator [Lepowsky, Frenkel and Merviman (1984)]. In the sense of Witten, this operator may be written as

$$V_{\nu\alpha} = \nabla_- x^L_\nu\, \nabla_+ x^R_\alpha \ , \tag{22}$$

where $\nabla_\pm = [(\partial/\partial_\tau) \pm (\partial/\partial_\sigma)]$, and x^L and x^R contain the left and right moving oscillations of the string. It follows then that $\big(\delta[x^\mu(\sigma,\tau) - q^\mu]$ has a Fourier transform $V(P^\mu(\sigma,\tau)] = \exp\,\{iP^\mu\, x_\mu(\sigma,\tau)\}\ \big)$,

$$\langle V_{30} \rangle = -\frac{1}{4}\, \delta(x_1)\, \delta(x_2) \ , \tag{23}$$

for $x^L_i = X!R_i = 0$ $(i = 1, 2)$ and for the position of the string $x_\alpha = x^L_\alpha + x^R_\alpha$.

Thus, from Eq. (21),

$$\Delta = -\frac{1}{8} g \int_S dx_1 \, dx_2 \, \delta(x_1) \, \delta(x_2) \neq 0 . \tag{24}$$

This shows that the superstring is an axion string and the boundary of an axion domain wall.

From this discussion, the author concludes that ϕ is an invisible axion, because ϕ has the standard coupling of an axion. Theories with axions usually have axion domain walls [Stecker and Shafi (1985), Vilenkin (1985)].

In resume, axions, whose role in phenomenology is well known [Peccei and Quinn (1977), Weinberg (1978), Wilczek (1978), Bardeen and Type (1978), Donnelly *et al.*, (1978), Georgi (1978), Dimopoulos and Susskind (1979), Kim (1979), Dicus *et al.* (1980), Chikashige *et al.* (1980), Dine *et al.* (1981), Sikivie (1982), Manohar (1982), Wacker (1983), Iwamoto (1984), Kaplan (1985), Nelson (1985), Cowan *et al.* / (1986), Tsi (1986), Bjorken *et al.* (1986), Pantziris and Kang (1986), Davis (1986), Dimopoulos *et al.* (1986)], are a general consequence [Witten (1984)] of the new anomaly cancellation mechanism [to the sense of Green and Schwarz (1984)].

We now discuss this new anomaly cancellation mechanism.

3.2 Position of the Problem: Axionic Anomaly

In general, by an anomaly we mean that some symmetry present in the classical action of a theory is not preserved by the full quantum theory. We will discuss some properties of the anomaly-free O(32) superstring theory recently discovered by Green and Schwarz (1984). It predicts, for example, axions in the model which are independent axions, in addition to the model dependent Peccei–Quinn symmetry. Axions appear through the zero modes of the antisymmetric tensor field B_{MN} (which is crucial in the anomaly cancellation mechanism), and their existence and properties are very much model independent.

We will discuss only the terms added to cancel the anomalies in the ten-dimensional theory, that is, to consider compactification from ten dimensions to $M^4 \times K$, M^4 being four-dimensional Minkowski space and K being a six-dimensional Calabi-Yan manifold, *i.e.,* a manifold of SU(3) holonomy. These terms to cancel the anomalies are important for the coupling of axions.

We consider the interaction of gravitation with the field of an antisymmetric tensor of rank $(p-1)$ and with a Yang–Mills vector field (such an interaction occurs in the boson part of supergravity and the effective superstring field theory) in the space M^d [Chapline and Mantou (1983), Derendiger *et al.* (1986)],

$$\begin{aligned} \hat{E}^{-1}\mathcal{L} \;=\; & -\frac{1}{2}\,R + \alpha\,\phi^{-3/2}\,H_{M_1\ldots M_P}\,H^{M_1\ldots M_P} \\ & + \beta\,\phi^{-3/4}\,F^{\sigma}_{MN}\,F^{\sigma MN} + \gamma\left(\frac{\partial_\mu \phi}{\phi}\right)^2 , \end{aligned} \tag{25}$$

where F is the curvature of the Yang–Mills field $(M, N, P = 1, \ldots, 10)$. $H = db$ is the field strength of a third rank antisymmetric tensor in eleven dimensions, $B_{M_1\ldots M_{P-1}}$. R is the scalar curvature, and ϕ is the scalar field. We consider values of α and β corresponding to both normal particles $(\alpha < 0, \beta < 0)$ as well as ghosts.

In order to obtain the $d = 4$ field theory from Eq. (25), one would have to integrate out the heavy modes in order to get an effective phenomenological Lagrangian. This is only possible using the following ansatz [Derendinger (1986), Sergré (1986)] for the metric

$$g_{MN} = \begin{pmatrix} e^{-3\gamma} g_{\mu\nu} & 0 \\ 0 & e^{\gamma} g_{mn} \end{pmatrix}, \tag{26}$$

where γ is a real scalar field, and g_{mn} is the internal metric tensor with unit determinant. The zehnbein determinant is then $\hat{E} = e^{-3\sigma} E$, where E is the four-dimensional vierbein determinant. This gives canonical Einstein Lagrangian in four space-time dimensions. The result [see Segré (1986)] is that among the gauge fields, the two dilatons ϕ and γ and their two pseudoscalar counterparts η and θ, arise from the antisymmetric tensor field $B_{\nu\nu}$. It is a pseudoscalar axion of four-dimensions coupling to $F \wedge F$. Further, the gravitino belongs to the four-dimensional representation of SO(6) $\simeq$ SU(4). The gauge fields require invariance under $\mathrm{SU(3)}_{\mathrm{holonomy}} \times \mathrm{SU(3)}_{\mathrm{gauge}}$. This leads to a single chiral 27–multiplet. More of the gauge fields mentioned above, we have graviton and gravitino, and $E_8 \times E_8'$ gauge fields, and gauginos.

One axion is model independent and arises from the mechanism of Witten (1985) which we have described in Sec. D.

3.3 Axion Coupling

Consider the following supergravity — super–Yang–Mills Lagrangian [Segré (1986)]:

$$\begin{aligned} \hat{E}\mathcal{L} \; = \; & -\frac{1}{2}\, R - \frac{1}{2}\, i\bar{\psi}_m \Gamma^{mnp} D_n \psi_p + \frac{i}{2}\, \bar{\lambda}\, \Gamma^m D_m \lambda \\ & + \frac{g}{16} \left(\frac{\partial_\mu \phi}{\phi} \right)^2 + \frac{3}{4}\, \phi^{-3/2} H_{mnp} H^{mnp} + \frac{3}{8}\, \sqrt{2}\, \bar{\psi}_m \frac{\partial\!\!\!/ \phi}{\phi}\, \Gamma^m \lambda \\ & - \frac{\sqrt{2}}{16}\, \phi^{-3/4} H_{mnp} \left(i\bar{\psi}_S\, \Gamma^{smnpr} \psi_r + 6 i \bar{\psi}^m \Gamma^n \psi^P + \sqrt{2}\, \psi_S\, \Gamma^{mnp} \Gamma^s \lambda \right) \end{aligned}$$

where λ is the right-handed spinor field ψ_{11}^{R} [$\psi_{11}^{R} \equiv \lambda$ by consistent truncation; $\psi_{11}^{L} \equiv (1+\gamma_{11})\,\psi_{11} = 0$, $Psi_{m}^{R} \equiv (1-\gamma_{11})\,\psi_{m} = 0$]. D is the covariant derivative, and Γ are the gamma matrices in eleven dimensions.

It is suggested by Witten (1985) and Segré (1986) that the term in Eq. (27) proportional to $H_{\mu}\nu\lambda\, H^{\mu}\nu\lambda$ can lead to an axion coupling proportional to $\theta\, F_{\mu}\nu\, \widetilde{F}^{\mu\nu}$. This comes from the product of $\partial_{\mu}\theta$ with the gauge field Chern–Simons term C^{μ} in H.

$$\partial_{\mu}\theta\, C^{\mu} \to \theta F \widetilde{F} \ , \tag{28}$$

that conduits to the classic axion coupling $\theta F\widetilde{F}$, where θ may be interpreted as the phase of a pseudoscalar field a whose coupling is of the form $(a/V_a)F_{\mu\nu}\widetilde{F}^{\mu\nu}$, where V_a is the scale of symmetry breaking. Thus, in order to have anomaly cancellation, additional terms [Green and Schwarz (1984), Alvarez–Gaumé and Witten (1984)] must be added to the Lagrangian, Eq. (25). One of these is of the form [Witten (1985), Segré (1986)]

$$\Delta\mathcal{L} \;=\; \mathcal{E}_{M_1\ldots M_{10}}\, B^{M_1M_2}\, F^{M_3M_4}\, F^{M_5M_6}\, F^{M_7M_8}\, F^{M_9M_{10}} = B\, Tr\, F^4 \ . \tag{29}$$

They arise, for example, for an anomaly of the form $(Tr\, F^2)(Tr\, F^4)$ in ten dimensions; that is, consider the H^2 term in the Lagrangian. Since $H = dB - W_{3Y} + W_{3L}$ where W_{3Y} is the Yang–Mills Chern–Simons three-form and W_{3L} is the Lorentz Chern–Simons three-form, has a w_{3y} form in it, H^2 may be written as

$$dB\, w_{3y} \sim B\, Tr\, F^2 \ , \tag{30}$$

where a massive B pole diagram leads to an effective $Tr(F^2)\, TR(F^4)$ interaction which cancels the anomaly [Green and Schwarz (1984)] in ten dimensions.

4. CONCLUSION AND DISCUSSION

In the minimal Standard Model the interaction in the Higgs sector conserve CP. This no longer holds if the Higgs sector is extended to contain more doublets.

Theoretically, the elementary Higgs sector of the standard electroweak model is unsatisfactory. Many ideas have been proposed to achieve a more fundamental understanding of the Higgs sector; prominant among the various ideas are technicolor and supersymmetry. Note that among all acceptable models of electroweak interactions only a supersymmetric extension of the Standard Model with two Higgs doublets is completely free of a strongly interacting Higgs sector. In all other cases, validity of a perturbation expansion at low energy imposes upper bounds on masses of Higgs bosons.

It is suggested by Wilczek (1978) and more recently by Eichten (1984) that the search for the Higgs boson in the decays of heavy quarkonium is a perfectly realistic experimental task.

I would like to stress that the low-energy theorems demonstrate a unique property of the Higgs bosons: if they can be observed experimentally, they make it possible to examine even smaller distances and to count the number of states with mass exceeding the mass of the scalar boson.

The situation might change if toponium [Franzini (1986), Gilman (1986)] with a mass $M_T < M_Z$ were to be discovered in the near future, or a search for radiative decays $T \rightarrow H^0 + \gamma$ were to be undertaken.

The axion has been intensively discussed. To all appearances, there are no particles with the properties predicted in the original papers.

Theories with axions usually have axion domain walls [Stecker and Shafi (1983)].

The invisible axion is an expected one in grand unified theories which have symmetry breaking scales of order $\sim 10^{15}$ GeV [Kim(1982)].

Axions are a general consequence [Witten (1984)] of the new anomaly cancellation mechanism [Green and Schwarz (1984)].

Axion strings are invariable superconducting [Callen and Harvey (1984), Lazarides and Shafi (1984), Rohm (1984)] with an asymmetry between the number of left-moving and right-moving charge carriers.

The superstring is an axion string, the boundary of an axion domain wall.

Besides solving the strong CP–problem, axions have been proposed as the missing mass of the universe.

It is suggested by Sikivie (1983) that the invisible axion might be observable experimentally from their conversion into γ–rays in a strong nonuniform magnetic field.

Efforts have been undertaken to detect these axions [Cowan *et al.* (1986), Bjorken *et al.* (1986), Tsai (1986), Pantziris and Kang (1986), Dearborn, Steigman and Schramm (1986)] using different methods.

Although the invisible axion was proposed to be a natural solution to the strong CP problem, its mass in the Dine, Fischler and Srednicki model is left undetermined. Since the coupling of this axion to matter is presumably very weak, laboratory experiments are unable to detect it and one can derive only bounds on its mass from astrophysical considerations (Table II resumes some recent results).

An important stimulus for this search is the recent experiment of the caps epos collaboration at the GSI at the Darmstadt (1986) [Cowan *et al.* (1986)] which observed anomalous pairs of electrons and positrons in the laboratory. It is probably a pseudoscalar boson of mass about 1.8 MeV. Theoretically, three new models [Peccei, Wu and Yanagida (1986), Krauss and Wilczek (1986), and Brodsky *et al.*] proposed accommodated the new data and the negative results.

NOTE ADDED TO THE PROOF: THE GAUGE GROUP PROBLEM The gauge group of the minimal Standard Model is SU(3) $\otimes$ SU(2) $\otimes$ U(1) that requires three independent coupling constants. One solution is the idea of super grand unification, *i.e.*, the assumption of the existence of a large simple group (or of a group which is a product of identical simple groups related by a discrete symmetry), in which the strong, electroweak and gravitational gauge groups are embedded. This reduces the number of independent gauge coupling constants to one. One such super (large) simple group [Ne'eman (1979, 1986)] could be F_1, the Monster group [Griess (1982)] which has been recently introduced in string theory [Chapline (1986)].

F_1 is a simple, finite group of Lie type (analogous to a Lie group). It is defined as automorphisms of a certain commutative nonassociative algebra of dimension 196884. F_1 presents:

- Connection with modular (invariant) forms $\longrightarrow$ string theories.
- Connection with vertex operators $\longrightarrow$ cosmic superstrings [Witten (1985)].
- Discrete symmetry $\longrightarrow$ supersymmetry.
- Complex irreducible representations (multiplets) $\longrightarrow$ electroweak theory.

- Connection with infinite dimensional Lie algebras, because of graded spaces where generating functions are modular forms $\longrightarrow$ superalgebra.
- Connection with Orbifolds [Bagger (1986)].

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FIGURE CAPTIONS

Fig. 1 A long superstring is running in the x_3 direction at rest ($\sigma = x_3'$ and $\tau = t$, time).

Fig. 2 (a) Axion production by the Compton process; (b) axion production by bremstrahlung (c) axion production by Primakoff process; (d) axion bremstrahlung production by lepton pair; (e) axion bremstrahlung production by an electron in the atomic targed [Tsai (1986), Cowan (1986)].

Fig. 3 Branching Ratio versus Higgs Mass: as these thresholds are crossed, the photon efficiency will most likely have discontinuities.

Fig. 4 Significance versus Photon Energy: standard derivations from smooth background.

Fig. 5 (a) $\mathrm{BR}(\Upsilon \to \gamma + X) \times 10^{-3}$ versus Photon Energy (90% C.L. upper limit); (b) $\mathrm{BR}(\Upsilon \to \gamma + X) \times 10^{-3}$ versus Higgs Mass (90% C.L. upper limit). Figure 4(b) indicates that only for Higgs mass around 5.5 GeV/c^2 does this analysis come close to the Wilczek (1977) estimate for the branching ratio of a minimal Higgs particle. For a Higgs mass below about 4 GeV/c^2, the efficiency drop, due to Bhabha rejection in the hadron selection routines, causes a large increase in the corresponding upper limit.

Fig. 6 The 90% C.L. upper limit of the BR($\Upsilon \to \gamma H^0$) versus M_{H^0}. The CUSB 90% C.L. upper limit of branching fraction for radiative Υ decay into minimal Higgs bosons.

Fig. 7 E_γ versus Recoil Mass:

$$\Upsilon(2s) \to \pi \circ \pi^\circ \, \Upsilon(1s)$$
$$\hookrightarrow \gamma\lambda\bar{\lambda}$$
$$(SUSY\,Goldstone fermions)$$

No events above $E_\gamma = 1200$ MeV in $\Upsilon(1s)$ bin.

TABLE I. **H I G G S**

(A) The resulting limits from PEP (MAC, MARK II, J), (CELLO, JADE), and TASSO [See Komamiya (1986)].

Year & Collaboration	Decay Mode & Branching Fraction	Remarks & Topology Studied
<u>1982</u> CELLO, JADE, MAC, MARK J, MARK II <u>1985</u> MARK J	$H^- \rightarrow \bar{\tau}\bar{\nu}_\tau$	The even shape is an acoplanar τ pair.
<u>1982</u> JADE, MAC, MARK J, MARK II <u>1985</u> MARK J	$H^- \rightarrow s\bar{c}$ or $b\bar{c}$	The even shape would be a τ accompanied by jets.
<u>1982</u> CELLO, JADE, MARK J, MARK II <u>1983</u> TASSO, CLEO <u>1985</u> MARK J	$BR(H^- \rightarrow \bar{\tau}\bar{\nu})$ + $BR(H^- \rightarrow \text{hadrons})$ $= 1$	$M_{H^+} < 5$ GeV (99.5% C.L.) The topology of the resultant limits from PEP and PETRA experiments are summarized in Konamiya (1986), *i.e.*, $BR(H^- \rightarrow \bar{\tau}\bar{\nu})$ *versus* $BR(H^- \rightarrow \textit{hadrons})$
<u>1983</u> TASSO	$H^\pm \rightarrow$ hadrons	The even shapes are four jets.
<u>1986</u> NE'EMAN	$M_H = 2M_W \sim 170$ GeV	Theory: using simple supergroup

TABLE I. **H I G G S.**

(B) Recent results from the CRYSTAL BALL at DORIS experiment [See Lowe (1986) and figures hère given.]

Year & Collaboration	Decay Mode & Branching Fraction	Remarks & Topology Studied
1986		Nondeviation from Standard Model
CRYSTAL BALL AT DORIS	$\Upsilon(1S) \to \gamma + x$	• Minimal Higgs sector
	Higgs Search $\left\{ \begin{array}{l} 2jets \\ \tau^+\tau^- \end{array} \right.$	$f \searrow$ …… H, e.g., $e^+e^- \to H$ $\nearrow f$
	Axion (unseen)	This coupling makes production difficult.
	(Nonres.) (unseen)	• $BR[\Upsilon \to \gamma H] \simeq 2.3 \times 10^{-4} \times (1 - M_H^2/M_\Upsilon^2)$ (minimal Higgs)
		• $H \to \tau\bar{\tau}$, $c\bar{c}$: $M_h \gtrsim 3710$ MeV
		• $\Upsilon(1S) \to \gamma +$ (unseen): M (unseen) O (GeV)

1) Search for $\Upsilon(1S) \to \gamma +$ Higgs

- Experimentally $\Upsilon(1S) \to \gamma+$ Higgs search just starting to test theory. $BR[\Upsilon \to \gamma + H] \simeq 2.3 \times 10^{-4}\ [1 - (M_H^2/m_\Upsilon^2)]$ (minimal Higgs).
- Lower limit O (few GeV) quite possible but Weinberg–Linde limit (~ 7.3 GeV) will be very difficult. Theory may drop predicted rates again.
- Radiative Υ decays very difficult — none seeen yet.

2) Search for $\Upsilon(1S) \to \gamma + H$
$$\hookrightarrow \tau^+\tau^-$$

- In minimal version of Higgs sector $H \to \left. \begin{array}{ll} c\bar{c} & 72\% \\ \tau\bar{\tau} & 22\% \end{array} \right\} M_H \gtrsim 3730 \text{ MeV } (c\bar{c})$

  Thus $\Upsilon \to \gamma + X$ is more sensitive to Higgs production.
- In some nonminimal models $H \to c\bar{c}$ suppressed and $H \to \tau\bar{\tau}$ decay dominant.

3) Search for $\Upsilon(1S) \to \gamma+$ unseen.

- Sensitive measurement to $E_\gamma < E_{beam}$. Thus M (unseen) $\sim$ O(GeV).
- No events above $E_\gamma = 1200$ MeV in $\Upsilon(1S)$ bin. Thus, $BR[\Upsilon(1S) \to \gamma +$ unseen$]$ $< 2.3 \times 10^{-3}$ for $\tau > 10^{-7}$ sec $0 < M$ (unseen) < 8.1 GeV . (See Fig. 6.)

TABLE II. **A X I O N**

Astrophysical bound s on the axion mass. $\widetilde{V}_a$ is the mass scale for the Peccei–Quinn symmetry violation (10^7 GeV $\lesssim \widetilde{V}_a \leq 10^9$ GeV).

Year	Group	m_a	$\tilde{V}_a$	Production of Axion by
1975	Sato and Sato	> 0.35 MeV	No bound	Super giant stars
1978	Discus *et al.*	> 0.2 MeV	No bound	Red Giant Stars
1980	Dicus *et al.*	< 0.01 eV	$> 10^9$ GeV	Red Giant Stars
1982	Fukugita *et al.*	< 1 eV	$> 4 \times 10^7$ GeV	
1982	Watamura and Yoshimura	$\sim 1.3 \times 10^{-14}$ MeV	$\gtrsim 10^{15}$ GeV	Higgs Decay
1984	Iwamoto	$\sim 3 \times 10^{-8}$ eV	$> 3 \times 10^9$ GeV	Neutron stars
1985	Kaplan	< 23 eV	No bound	$a\gamma\gamma$
1986	Cowan[a] *et al.*	~ 1.8 MeV	No bound	Electron beam
1986	Tsai[a]	~ 1.7 MeV	$\sim 10^2$ GeV	Proton beam, photon beam
1986	Pantziris *et al.*	< 2.7 eV	$> 10^7$ GeV	Sun

[a]Laboratory experiments.

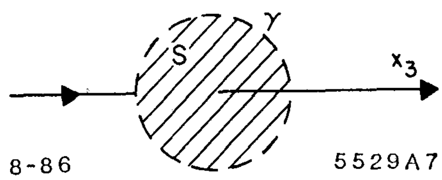


Fig. 1

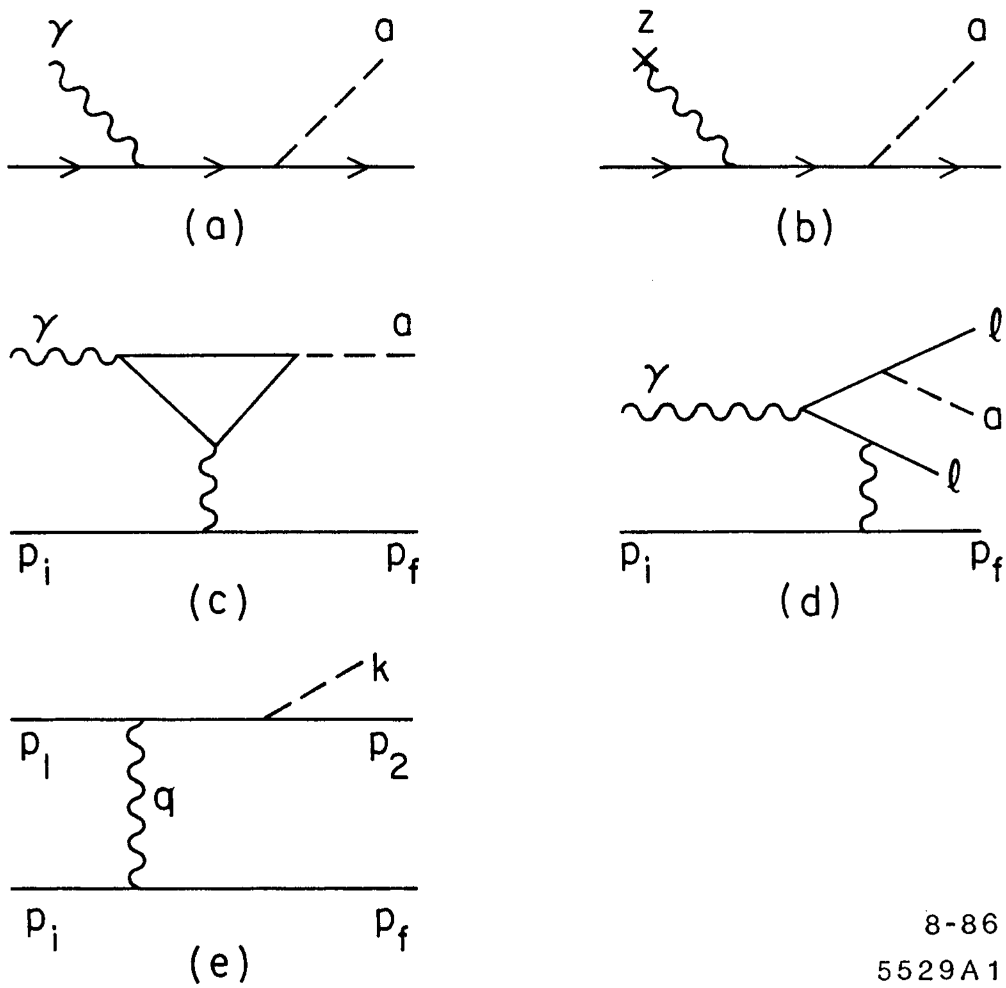
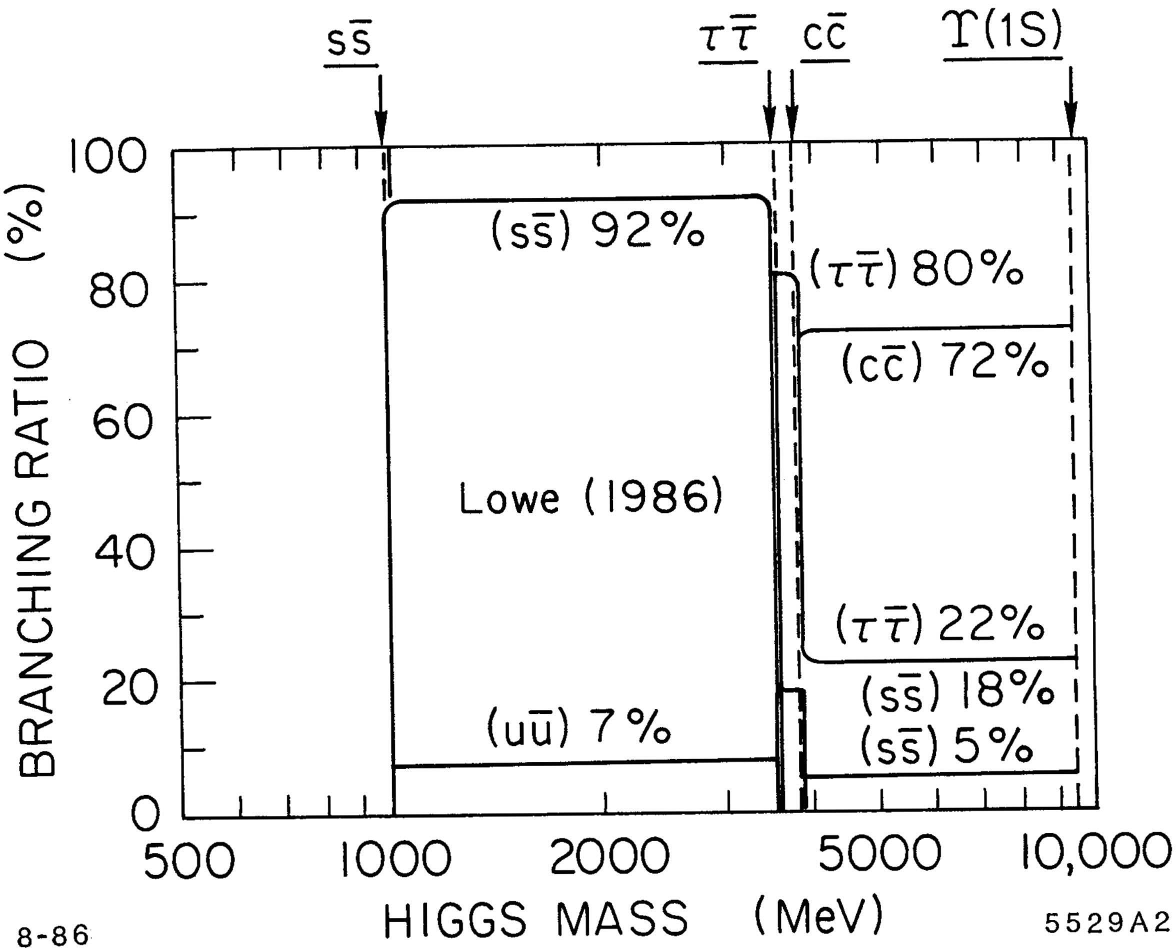
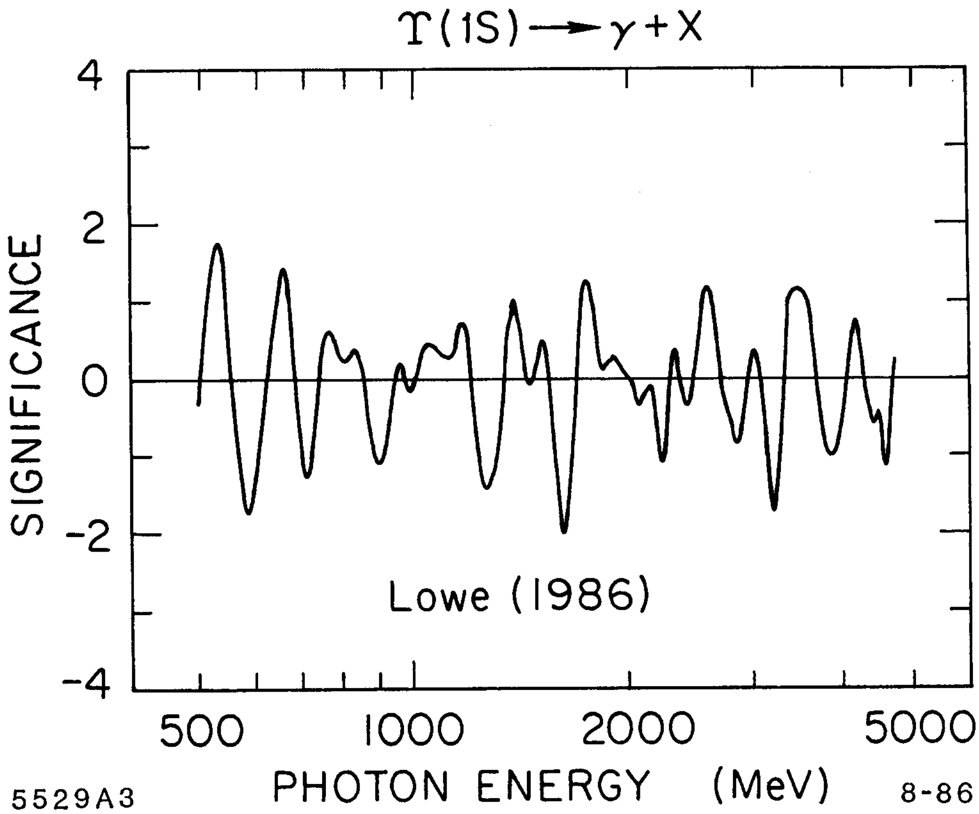


Fig. 2



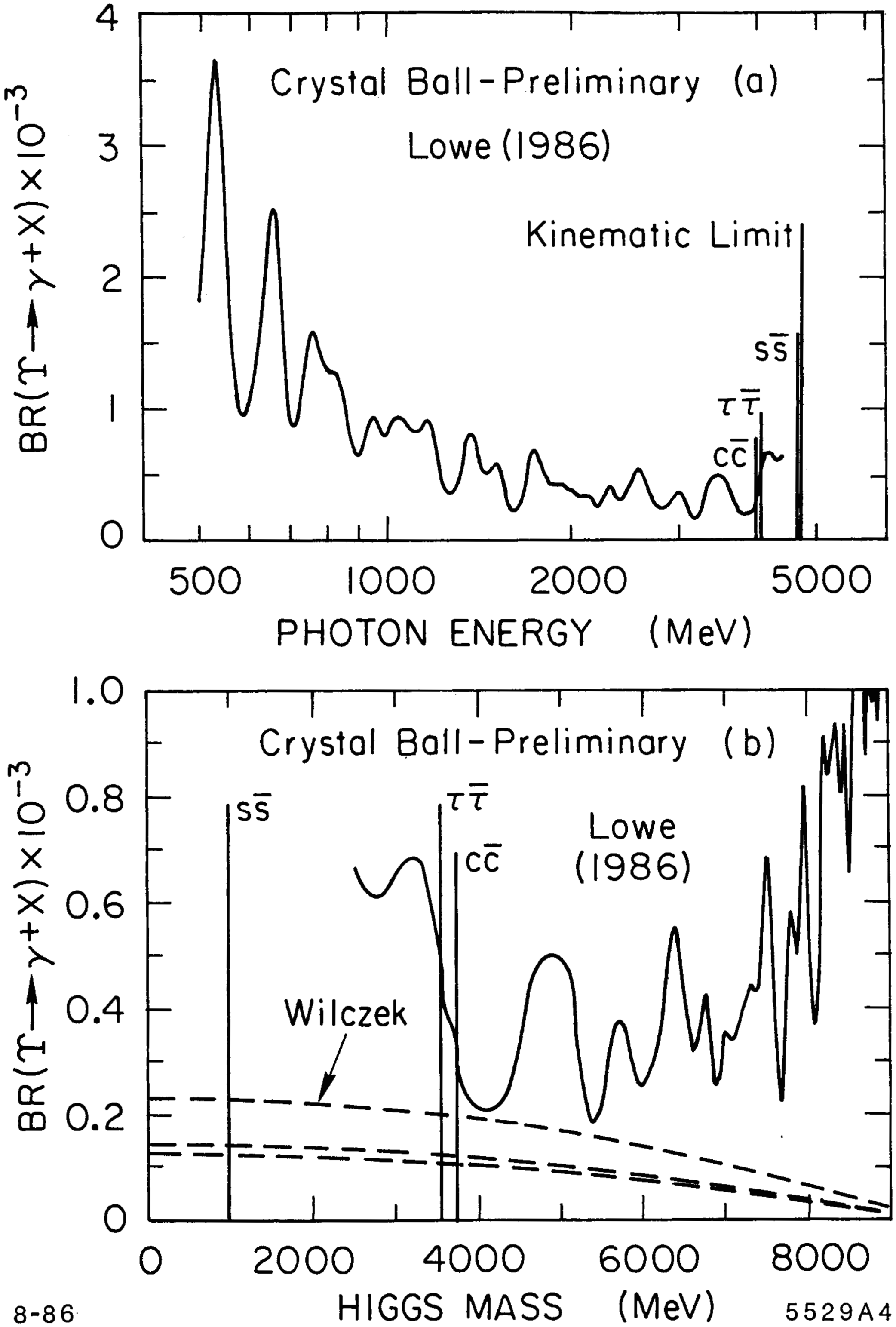
$s\bar{s}$
$\tau\bar{\tau}$
$c\bar{c}$
$\Upsilon(1S)$
100
80
60
40
20
0
BRANCHING RATIO (%)
$(s\bar{s})$ 92%
$(\tau\bar{\tau})$ 80%
$(c\bar{c})$ 72%
Lowe (1986)
$(\tau\bar{\tau})$ 22%
$(s\bar{s})$ 18%
$(u\bar{u})$ 7%
$(s\bar{s})$ 5%
500
1000
2000
5000
10,000
HIGGS MASS (MeV)
8-86
5529A2

Fig. 3



$\Upsilon(1S) \longrightarrow \gamma + X$
4
2
0
-2
-4
SIGNIFICANCE
Lowe (1986)
500
1000
2000
5000
PHOTON ENERGY (MeV)
5529A3
8-86

Fig. 4



Crystal Ball-Preliminary (a)
Lowe (1986)
Kinematic Limit
$s\bar{s}$
$\tau\bar{\tau}$
$c\bar{c}$
BR($\Upsilon \to \gamma + X$) $\times 10^{-3}$
PHOTON ENERGY (MeV)
Crystal Ball-Preliminary (b)
$s\bar{s}$
$\tau\bar{\tau}$
$c\bar{c}$
Lowe (1986)
Wilczek
BR($\Upsilon \to \gamma + X$) $\times 10^{-3}$
HIGGS MASS (MeV)
8-86
5529A4

Fig. 5

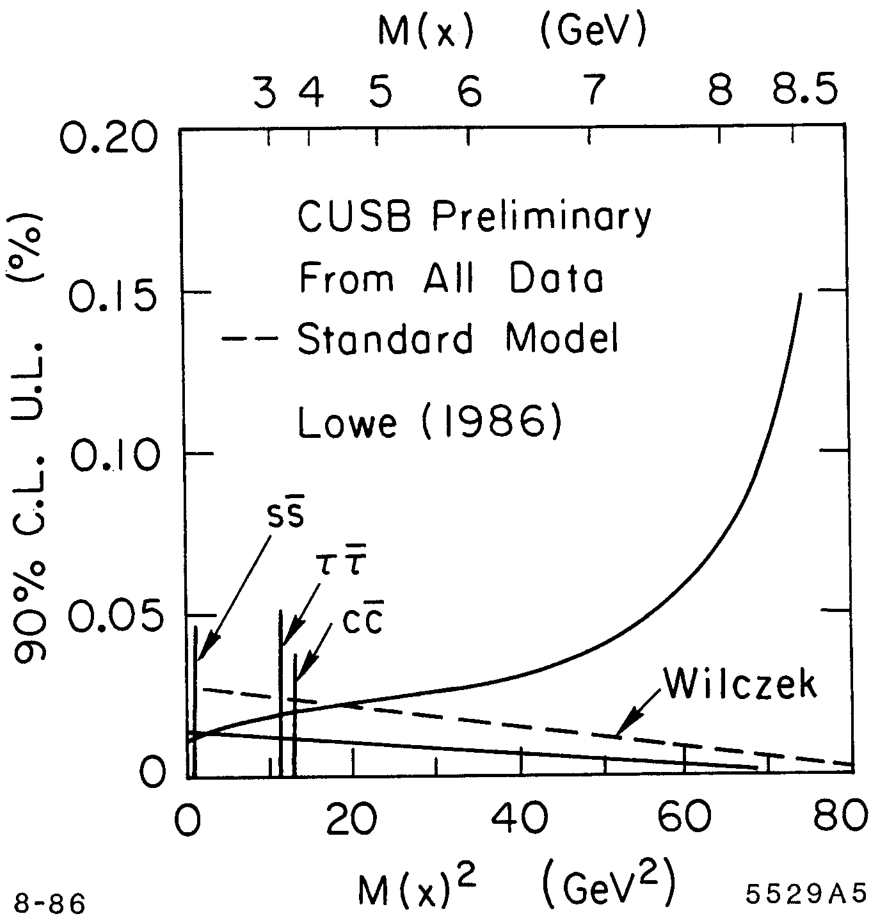
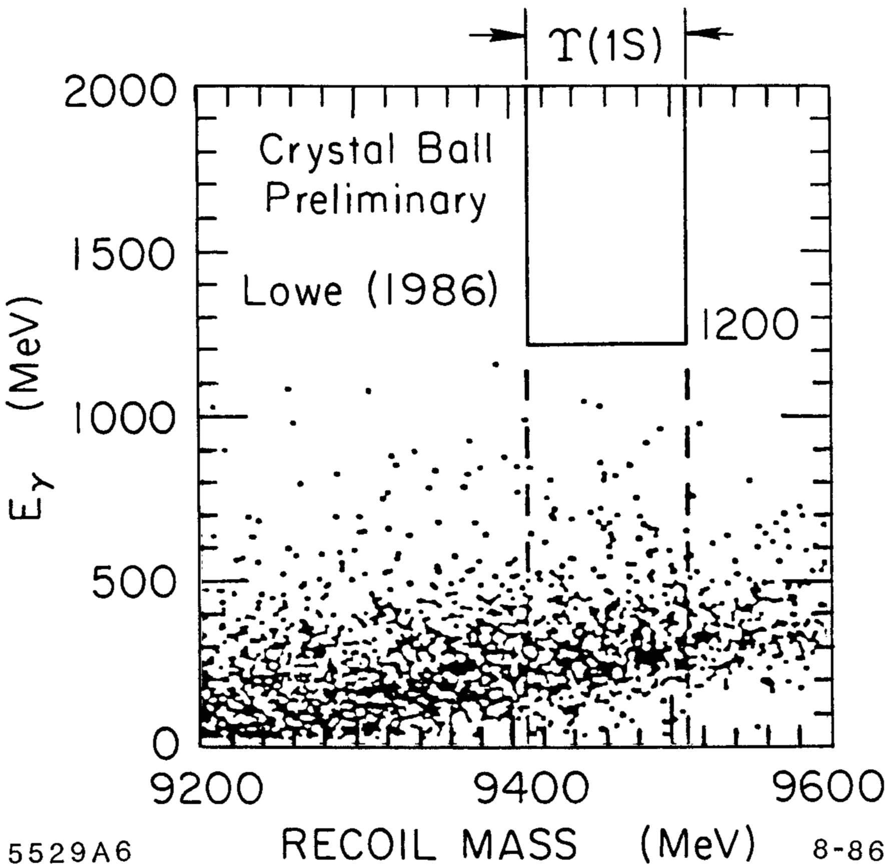

M(x) (GeV)
3 4 5 6 7 8 8.5
0.20
0.15
0.10
0.05
0
90% C.L. U.L. (%)
CUSB Preliminary
From All Data
-- Standard Model
Lowe (1986)
$s\bar{s}$
$\tau\bar{\tau}$
$c\bar{c}$
Wilczek
0 20 40 60 80
$M(x)^2$ (GeV^2)
8-86
5529A5

Fig. 6



Υ(1S)
2000
1500
1000
500
0
Crystal Ball
Preliminary
Lowe (1986)
1200
E_γ (MeV)
9200
9400
9600
RECOIL MASS (MeV)
5529A6
8-86

Fig. 7