Higgs particles

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Usp. Fiz. Nauk **145**, 185–223 (February 1985)

The theoretical work on models of the electroweak interaction and simple grand unified models with a nonstandard set of Higgs particles is reviewed. Emphasis is placed on light and even strictly massless Higgs particles: Goldstone and pseudo-Goldstone bosons. It is shown that such bosons arise in a natural way in the theory if the Higgs particles are in fact composite. The low-energy effective Lagrangian of these particles is studied. A detailed study is made of the problem of *CP* breaking in a strong interaction and of a natural solution of this problem through the introduction of a pseudo-Goldstone particle: an axion. The theory of the "standard" axion and its experimental status are reviewed. Possible "invisible" and "visualized" axions are discussed, as are certain astrophysical aspects of the existence of an axion. By analogy with the axion, an analysis is made of another hypothetical particle: the strictly massless Goldstone boson or arion. Model-independent properties of the arion are determined. The similarity between the arion fields and magnetic fields and the differences between these fields are shown. Possible methods for detecting an arion field are discussed. An experiment which has set a limit on the strength of the arion interaction is described. Neutral Goldstone bosons whose emission is accompanied by changes in fermion flavors ("familons") are discussed. Two versions of the theory with a Goldstone boson (a majoron) which arises upon a spontaneous breaking of lepton number are described.

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INTRODUCTION

The successful discovery of the W and Z bosons on the colliding-beam accelerator at CERN[1,2] erased essentially all doubt regarding the validity of that part of the Glashow-Weinberg-Salam unified theory of the electroweak interaction[3] which pertains to the interaction of vector bosons with quarks and leptons. Actually, there had been no particular doubt regarding the validity of the description of the interaction of these particles at low energies since the mid-1970s, when the existence of neutral currents was verified.[4] The discovery of the carriers of the electroweak interaction—the vector bosons—was the triumphant culmination of the first stage of the verification of the theory.

The next and perhaps decisive step in the verification of the entire concept underlying the theory would appear to be the discovery of Higgs bosons. Indeed, the most important property of the Glashow-Weinberg-Salam theory is its renormalizability. No one has yet been able to construct a renormalizable theory containing massive vector bosons without introducing spinless fields. Since Higgs bosons are required for renormalizability and, in this sense, for self-consistency of the theory, a search for these bosons would appear to be one of the most important tasks for the immediate future in the physics of elementary particles.

In the simplest version of the theory there is only one elementary Higgs doublet. This circumstance corresponds to a situation in which, after the Goldstone degrees of freedom have been used up in giving mass to the charged W bosons and the neutral Z boson, only a single observable neutral Higgs boson remains (see Ref. 5, for example). A large part of the present review will be a description of the properties of this "standard" Higgs boson.

However, there is no special reason to assume that the Higgs sector contains only a single boson. The fact that we have so far not observed a single Higgs boson is not evidence that Higgs particles of only a single type exist. Furthermore, there are several theoretical factors which make it desirable to expand the Higgs sector. We will take up these factors in Sec. 3, where we will discuss "nonstandard" Higgs particles.

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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.

Here is a brief outline of the present review. In Sec. 1 we briefly list the properties of the standard Higgs boson: We describe its interaction with quarks, leptons, and gauge bosons; the basic decay modes; and theoretical limitations on its mass. In Sec. 2 we describe reactions in which a neutral Higgs boson could be produced. In particular, we discuss in detail what appears to be the most promising reaction: that involving the joint production of a Higgs boson with gauge bosons. In Sec. 3 we discuss nonstandard Higgs bosons: charged bosons and light scalar particles which arise in models with dynamic symmetry breaking (the axion, the arion, the familons, and the majoron). In Sec. 4 we discuss the present experimental status of the standard Higgs boson and of charged bosons. Section 5 deals with the outlook for the discovery of Higgs bosons on the accelerators which are expected to come on line in the near future.

We will not discuss many of the theoretical aspects of the existence of Higgs bosons. Some of these aspects have been reviewed by Vaĭnshteĭn *et al.*[6] The possibilities of a search for Higgs bosons were first discussed by Bogomol'nyĭ[7] and Ellis *et al.*[8] The exceptional role which Higgs bosons should play in a test of the concept of a spontaneously broken gauge invariance was emphasized in a paper by Okun' at a Bonn conference.[9] Among other reviews of Higgs particles we might mention the well-known review by Gaillard[10] and the lectures by Ellis.[11] The possibilities of searching for Higgs bosons on accelerators at high energies were discussed in the reviews by Barbiellini *et al.*[12] and Ali.[13] Finally, we note that certain questions relating to the properties of Higgs bosons were discussed in detail in a lecture by the present authors at a school of the Leningrad Institute of Nuclear Physics.[14]

1. THE STANDARD HIGGS BOSON AND ITS EXPECTED PROPERTIES

The minimal version of the $SU(2)_L \times U(1)_Y$ theory of the electroweak interaction requires the existence of only one Higgs multiplet: the doublet $\varphi = \binom{\varphi^+}{\varphi_0}$ with hypercharge $Y = 1$. In the case of a spontaneous development of a nonvanishing vacuum expectation value of the field φ

$$\varphi(x) = \begin{pmatrix} \varphi^+(x) \\ \varphi^0(x) \end{pmatrix} = \begin{pmatrix} \varphi^+(x) \\ \frac{1}{\sqrt{2}}(v + H(x) + i\chi(x)) \end{pmatrix}, \quad v \neq 0, \qquad (1.1)$$
$$\langle H(x)\rangle = \langle \chi(x)\rangle = 0,$$

the $SU(2)\times U(1)$ gauge group is broken down to the $U(1)_{em}$ group. The W and Z bosons and the fermions acquire masses

$$m_W^2 = g^2\frac{v^2}{4}, \quad m_Z^2 = \bar{g}^2\frac{v^2}{4}, \quad m_f = h_f\frac{v}{\sqrt{2}}; \qquad (1.2)$$

here $\bar{g}^2 = g^2 + g'^2$ and h_f are Yukawa constants, and the scalar fields $\varphi^+(x)$ and $\chi(x)$ disappear from the spectrum of physical states by virtue of the Higgs mechanism. Only the single neutral scalar boson H^0 turns out to be observable.

In a minimal scheme with a single Higgs doublet the vacuum expectation value v can be expressed unambiguously in terms of the Fermi constant G_F with the help of (1.2):

$$v = (G_F\sqrt{2})^{-1/2} \approx 246\,\text{GeV}. \qquad (1.3)$$

The self-effect of the Higgs field and all the interactions of the Higgs boson H^0 with gauge bosons and fermions are determined entirely by the masses of these particles:

$$\begin{aligned} L_H = {}& \frac{1}{2}(\partial_\mu H)^2 - \frac{m_H^2}{2}H^2 - \frac{m_H^2}{2v}H^3 \\ & - \frac{m_H^2}{8v^2}H^4 + \frac{2m_W^2}{v}W_\mu^+W_\mu^-H \\ & + \frac{m_Z^2}{v}Z_\mu^2H + \frac{m_W^2}{v^2}W_\mu^+W_\mu^-H^2 + \frac{m_Z^2}{2v^2}Z_\mu^2H^2 - \frac{m_f}{v}\bar{f}fH. \end{aligned} \qquad (1.4)$$

Using (1.2), we can write the interaction of H with the W and Z bosons in the more customary form

$$\frac{g^2v}{2}W_\mu^+W_\mu^-H + \frac{g^2}{4}W_\mu^+W_\mu^-H^2 + \frac{\bar{g}^2v}{4}Z_\mu^2H + \frac{\bar{g}^2}{8}Z_\mu^2H^2.$$

Lagrangian (1.4) reflects a basic property of Higgs bosons: Their interaction with particles is proportional to the mass of these particles (in the case of fermions, the mass appears in the amplitude, while in the case of bosons the square of the mass appears in the amplitude). The appropriate dimensionality of the corresponding terms in the Lagrangian is ensured by the powers of the vacuum expectation value v.

Although the Higgs boson has no direct interaction with gluons or photons, it does appear on a single-loop level. The distinctive feature of the interaction of H^0 with fermions and vector bosons which we mentioned earlier has the consequence that the contribution of virtual heavy particles is independent of their mass for the amplitudes for the transitions $H^0 \to gg$ (g is the gluon) and $H^0 \to \gamma\gamma$. The effective Lagrangian for the interaction with gluons, for example, is[15,16]

$$L_{eff} = \frac{n_h}{v}\frac{\alpha_s}{12\pi}G_{\mu\nu}^aG_{\mu\nu}^aH, \qquad (1.5a)$$

while that for the interaction with photons is[8,17,18]

$$L_{eff} = \frac{\alpha}{8\pi v}\left(-7 + \frac{4}{3}\sum_f Q_f^2\right)F_{\mu\nu}F_{\mu\nu}H. \qquad (1.5b)$$

Here n_h is the number of heavy quarks, and the sum over f is carried out over all the heavy fermions. More-regorous expressions, which incorporate intermediate particles of arbitrary masses, are given in Refs. 5, 17, and 18, for example. Figure 1 shows a plot of the fermion contribution as a function of the ratio m_f/m_H. We see that the form factor is indeed approximately unity at $m_f \gtrsim 0.2\, m_H$. At small values of

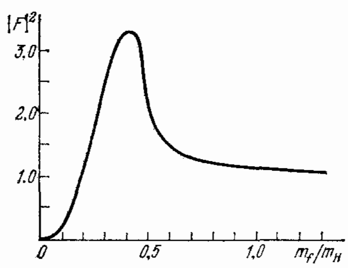

FIG. 1. The form factor F of the $H^0 \to gg$ ($H^0 \to \gamma\gamma$) transition as a function of the ratio m_f/m_H.

m_f/m_H, the form factor falls off rapidly since it contains a factor of m_f^2/m_H^2.

a) Mass of a Higgs boson

The mass of a Higgs boson is the only adjustable parameter in Lagrangian (1.4). This mass is related unambiguously to the self-effect constant of the Higgs field, λ:

$$\frac{\lambda}{4!}H^4=\frac{m_H^2}{8v^2}H^4. \tag{1.6}$$

Consequently, limitations on m_H can be found in terms of permissible values of λ.

If we are to be able to apply a perturbation theory to the Higgs particles, the constant λ must be quite small. It is reasonable to assume that λ cannot exceed a critical value at which the Born amplitude for the scattering $H^0+H^0 \to H^0+H^0$ reaches the unitary limit. The Born amplitude (normalized in such a way that its square is equal to the cross section) for interaction (1.6) is $f=-\lambda/8\pi\sqrt{s}$, where $\sqrt{s}$ is the energy. Since the scattering occurs in the S wave, it is a simple matter to find a limitation on λ corresponding to the unitary limit $|f|<1/k$ (k is the momentum; at high energies we have $k\approx\sqrt{s}/2$): $\lambda<16\pi$. Hence

$$m_H^2\leqslant\frac{8\pi\sqrt{2}}{3G_F}=(1\,\mathrm{GeV})^2. \tag{1.7}$$

A more rigorous discussion is based on a study of the three-channel condition for unitarity for a system of longitudinally polarized W^+W^-, Z^0Z^0, and a pair of Higgs bosons, H^0H^0 (Refs. 19 and 20). The idea behind this approach is that the scattering amplitude of longitudinally polarized W and Z bosons, without allowance for the exchange of a Higgs boson, increases quadratically with their energy, and an unacceptable growth of the cross sections is stopped at $s\gtrsim m_H^2$ only when the Higgs boson comes into play (see Ref. 5, for example). It is thus clear that m_H cannot be arbitrarily large, for otherwise the scattering amplitudes would increase beyond the unitary limit. The restriction on the mass of the H boson found by this approach is exactly the same as (1.7).

Incidentally, an upper limit on m_H should not be understood literally as the necessary existence of an elementary scalar Higgs boson with a mass less than 1 TeV. Instead of this boson, the theory might contain, for example, composite scalar bosons with a mass $\lesssim 1$TeV.

If we are to have a relatively light Higgs boson, the self-effect constant of the Higgs field must be small: $m_H^2=(\lambda/3)v^2$. The self-effect of Higgs fields cannot, however, be arbitrarily small: Even if their nucleating self-effect were not present, the quantum interaction with W and Z bosons would lead to a value $\lambda\sim g^4$, $\bar{g}^4$.

The complete effective potential of the Higgs doublet field, φ, (1.1), takes the following form when single-loop corrections are taken into account[21]:

$$V_{\mathrm{eff}}=\mu^2|\varphi|^2+\frac{1}{64\pi^2}\Big(3\sum_V m_V^4(\varphi)+\sum_S m_S^4(\varphi)-4\sum_f m_f^4(\varphi)\Big)\ln\frac{|\varphi|^2}{M^2}, \tag{1.8}$$

where $m_V(\varphi)$, $m_S(\varphi)$, and $m_f(\varphi)$ are the masses of vector bosons, scalar particles, and fermions in the external field φ; μ^2 and M^2 are arbitrary normalization constants; and the summation in (1.8) is over all the particles with the given spin. In the standard model we would have

$$m_W^2(\varphi)=\frac{g^2}{2}|\varphi|^2,\quad m_Z^2(\varphi)=\frac{\bar{g}^2}{2}|\varphi|^2,$$

$$m_\gamma^2(\varphi)=0,\quad m_f(\varphi)=h_f|\varphi|,$$

where h_f is the Yukawa constant, given by $h_f=m_f\sqrt{2}/v$, where m_f is the physical mass of the fermion. As for the mass of the scalar boson, m_S^2, we note that since we are interested in relatively light scalar bosons we ignore their contribution to V_{eff} in comparison with that of gauge bosons [as will be seen below, the lower boundary on m_H^2 is, in order of magnitude, $g^2M_W^2\sim(7\,\mathrm{GeV})^2$; i.e., the contribution of scalar particles can in fact be ignored].

The potential $V_{\mathrm{eff}}(\varphi)$ must have a minimum at $|\varphi|=\langle\varphi^0\rangle=v/\sqrt{2}$. This condition establishes a relationship between M^2 and μ^2:

$$\gamma\left(\ln\frac{v^2}{2M^2}+\frac{1}{2}\right)=-\frac{\mu^2}{v^2},\quad \gamma=\frac{3m_Z^4+6m_W^4-4\sum_f m_f^4}{16\pi^2v^4}>0. \tag{1.9}$$

As a result, V_{eff} depends on only a single unknown parameter, say μ^2:

$$V_{\mathrm{eff}}(\varphi)=\mu^2|\varphi|^2+\gamma|\varphi|^4\left(\ln\frac{2|\varphi|^2}{v^2}-\frac{1}{2}\right)-\frac{\mu^2}{v^2}|\varphi|^4. \tag{1.10}$$

If $\mu^2<0$ (curves 1 and 2 in Fig. 2), the point $\varphi=0$ corresponds to a local maximum of the potential, and the ground state corresponds to $|\varphi|^2=v^2/2$. If $\mu^2>0$, there is, in addition to the minimum at $|\varphi|^2=v^2/2$, a minimum of V_{eff} at $\varphi=0$ (curves 3 and 4). This feature corresponds to a phase with unbroken symmetry. If $\mu^2\geqslant\gamma v^2$, the minimum at the point $|\varphi|^2=v^2/2$ disappears (curve 5). If the true vacuum is to correspond to a spontaneously broken symmetry, the following conditions are necessary:

$$V_{\mathrm{eff}}(\varphi)|_{|\varphi^2|=v^2/2}<V_{\mathrm{eff}}(\varphi=0)=0,\quad\text{or}\quad 2\mu^2<\gamma v^2.$$

The mass of a Higgs boson is determined by the second derivative of the potential with respect to the Higgs field at the point of the physical vacuum:

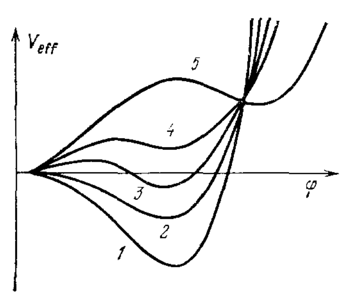

FIG. 2. The effective potential V_{eff} for various values of the parameter μ^2.

$$m_{\rm H}^2 = \frac{\partial^2 V_{\rm eff}}{\partial H^2} = \frac{1}{2}\,\frac{\partial^2 V_{\rm eff}}{\partial\,|\,\varphi\,|^2}\bigg|_{\varphi=\langle\varphi\rangle} = 2\gamma v^2 - 2\mu^2 = \gamma v^2 + (\gamma v^2 - 2\mu^2). \tag{1.11}$$

Accordingly, if our vacuum with a broken SU(2)×U(1) symmetry is in fact the state with the lowest energy, the following condition must hold:

$$m_{\rm H}^2 > \gamma v^2 = \frac{3m_{\rm Z}^4 + 6m_{\rm W}^4 - 4\sum_f m_f^4}{16\pi^2 v^2}. \tag{1.12}$$

The masses of the quarks and leptons which have been discovered to date are substantially smaller than the masses of the W and Z bosons. If we also ignore the t quark in comparison with $m_{\rm Z,W}$, condition (1.12) becomes

$$m_{\rm H} \geqslant \frac{\alpha}{4\sin\theta_{\rm W}}\sqrt{3[2+(\cos\theta_{\rm W})^{-4}]}\,v \approx 6.5\ \text{GeV}$$
$$\text{at}^{1)}\ \sin^2\theta_{\rm w} = 0.23. \tag{1.13}$$

This inequality is known as the Linde-Weinberg limitation.[23,24]

It should be stressed, however, that the limitation $m_{\rm H} \geqslant 6.5$ GeV applies only if the mass of the t quark is small. If m_t^4 cannot be ignored in comparison with the contribution of vector bosons in (1.12), the limitation on $m_{\rm H}$ becomes less stringent. When m_t approaches the critical value $m_0 = [(1/4)(m_Z^4 + 2m_w^4)]^{1/4} \approx 77$ GeV, the theory becomes unstable (as $m_t \to m_0$, the limitation on $m_{\rm H}$ diminishes to zero, and at $m_t > m_0$ we have $V_{\rm eff} \to -\infty$ as $|\varphi| \to \infty$). It might be concluded that in this theory there is a limitation on the mass of the t quark[25–29]; $m_t < 77$ GeV. This is true, however, only if the mass of the Higgs boson is not too large, specifically, if it is not comparable (in order of magnitude) with the mass of the vector bosons, in which case $m_{\rm H}^4$ would have to be added to the right side of (1.12). Clearly, the mass of the Higgs boson would have to be extremely large and able to stabilize the theory with $m_t > 77$ GeV (Fig. 3).

We thus see that the presence of a superheavy quark with $m_{\rm Q} < m_0$ lowers the Linde-Weinberg boundary (the "large top-quark-mass effect"):

$$m_{\rm H} > 6.5\ \text{GeV}\sqrt{1-\frac{m_{\rm Q}^4}{m_0^4}}, \quad m_0 \approx 77\,\text{GeV}. \tag{1.13a}$$

The joint limitations on the mass of the t quark and the Higgs boson were discussed in Ref. 30.

A case which is slightly special from the theory standpoint is that in which we have[21] $\mu^2 = 0$ in the renormalized potential $V_{\rm eff}(\varphi)$ in (1.8) and (1.10) (curve 2 in Fig. 2). This case corresponds to the possibility that the renormalized mass of the scalar particles in the phase with the unbroken SU(2)×U(1) symmetry will be zero. Here the mass of the Higgs boson will be about 9.2 GeV, $\sqrt{2}$ times the lower boundary [(1.13), (1.13a)], as can be seen from (1.11):

$$m_{\rm H}^{\rm CW} \approx \frac{\alpha}{2\sqrt{2}\sin^2\theta_{\rm W}}\sqrt{3\left(2+\frac{1}{\cos^4\theta_{\rm W}}\right)}\sqrt{1-\frac{m_t^4}{m_0^4}}\,v \approx 9.2\ \text{GeV}. \tag{1.14}$$

The Linde-Weinberg limitation (1.13) may be violated if we assume that the existing physical vacuum is a metastable state. If, in the Glashow-Weinberg-Salam model with a single Higgs doublet, it turned out that the mass of the Higgs boson was less than 6.5 GeV (and this result was not due to heavy fermions), then we would sooner or later go into a state with $\langle\varphi\rangle = 0$. The lifetime of a metastable state of this sort, however, might be very long. According to Linde's estimates,[31] for Higgs-boson masses above 260 MeV the average time for a transition to the ground state would be more than 10^{10} yr. Higgs bosons with an extrememly small mass, $m_{\rm H} \geqslant 260$ MeV, might be possible from this point of view.

Analysis of the very early stages in the evolution of the universe, when the temperature was of the order of $T \sim 10$–10^3 GeV, imposes a considerably more severe limitation on $m_{\rm H}$. During the initial stage of the evolution, at a higher temperature, the effective potential had a unique minimum at the point $\varphi = 0$, and the symmetry was restored. If $m_{\rm H}$ is less then $m_{\rm H}^{\rm CW}$ in (1.14), then even at $T = 0$ there would be a local minimum at $\varphi = 0$, although this minimum would not be as deep as that at $|\varphi| = v/\sqrt{2}$ (Fig. 2). Analysis of the kinetics of the corresponding phase transition has shown[30,32,33] that if $m_{\rm H}$ is even 1% less then $m_{\rm H}^{\rm CW}$ in (1.14) a transition to our vacuum could not have occurred over the entire time the universe has existed, $t \sim 10^{10}$ yr, and after the phase transition the universe would have ended up unacceptably inhomogeneous and anisotropic.

In summary, the theoretical limitations on the mass of the Higgs boson which we have at the moment are extremely weak. Specifically, if superheavy fermions are ignored, they are

1 TeV $\gtrsim m_{\rm H} \gtrsim$ 6.5 GeV for a stable vacuum,
$\gtrsim$ 260 MeV for a metastable vacuum,
$\gtrsim$ 9.2 GeV when cosmological aspects of the evolution of the universe are taken into acount.

The complete set of limitations on the masses of fermions and the Higgs boson is given in Fig. 3, taken from Ref. 30.

b) Basic decay modes and lifetime

Since the coupling of the Higgs boson with other particles is proportional to the mass of these other particles, the Higgs boson decays primarily into quarks and leptons with the maximum possible mass. Consequently, both the lifetime and the relative probabilities of the specific decays modes depend strongly on the mass of the H^0 boson.

The mass interval 0.3–1 GeV is dominated by the decays $H^0 \to \mu^+\mu^-$ and $H^0 \to \pi\pi$:

$$\Gamma(H^0 \to \mu^+\mu^-) = \frac{G_F m_{\rm H} m_\mu^2}{4\sqrt{2}\pi}\left(1-\frac{4m_\mu^2}{m_{\rm H}^2}\right)^{3/2} \approx 7\ \text{eV}. \tag{1.15}$$

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1) In these estimates we are using the effective experimental value $\sin^2\theta_{\rm W} = 0.23$, ignoring the electroweak radiative corrections (see Ref. 22, for example).

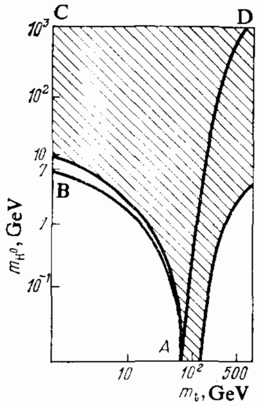

FIG. 3. The hatched region corresponds to those masses of the Higgs boson (m_H) and the t quark [more precisely, $(\Sigma_q m_q^4 + \frac{1}{3}\Sigma_l m_l^4)^{1/4}$] which are allowed by cosmological considerations. The region inside curve ABCD is the region of absolute stability of the phase with spontaneously broken symmetry (the limitations correspond to Ref. 30).

As for the decay into two pions, we note that it is substantially stronger than the naive estimate $\Gamma(H^0 \to \pi\pi) \approx \Gamma(H^0 \to u\bar{u}, d\bar{d})$. Specifically,[6]

$$\frac{\Gamma(H^0 \to \pi^+\pi^-)}{\Gamma(H^0 \to \mu^+\mu^-)} \approx \frac{m_\pi^4}{2m_\mu^2 m_H^2} |f(m_H^2)|^2 \frac{|\mathbf{p}_\pi|}{|\mathbf{p}_\mu|}, \qquad (1.16)$$

where f is the scalar form factor of the π meson [in the chiral limit we would have $f(0) = 1$].

If $m_H > 1$ GeV, the decays into mesons containing s quarks become the basic decay modes. For an even heavier H^0, the channels $H^0 \to \tau^+\tau^-$, $H^0 \to c\bar{c}$, $H^0 \to b\bar{b}$, etc., become dominant (Table I and Fig. 4). If the mass of the H^0 boson is above the threshold for the production of a pair of W or Z bosons, the decay modes $H^0 \to W^+W^-$ and $H^0 \to Z^0Z^0$ are dominant (if there is no heavy quark or lepton with a mass greater than $M_{Z,W}$)[19]:

$$\Gamma(H^0 \to W^+W^-) = \Gamma_0 (1-x)^{1/2} \left(1 - x + \frac{3}{4} x^2\right), \qquad (1.17)$$

$$\Gamma(H^0 \to Z^0Z^0) = \frac{1}{2} \Gamma_0 (1-y)^{1/2} \left(1 - y + \frac{3}{4} y^2\right),$$

$$\Gamma_0 = \frac{G_F m_H^3}{8\sqrt{2}\pi} \approx 41\,\text{GeV} \cdot \left(\frac{m_H}{500\,\text{GeV}}\right)^3,$$

$$x = \frac{4m_W^2}{m_H^2}, \quad y = \frac{4m_Z^2}{m_H^2}.$$

Since such decay modes have clear experimental manifestations, they are presently the leading candidates in a search for Higgs bosons with masses up to 600–700 GeV (Sec. 2). An even heavier H^0 boson would apparently be essentially impossible to observe directly because of the rapid increase in its width.

We can also write expressions for the widths of the decay of H^0 into two photons and two gluons[6]:

TABLE I.

m_H	Basic decay modes	Width
0–1.1 MeV	$H^0 \to 2\gamma$	$0.1\,\text{eV} \cdot \left(7 - \frac{4}{3} \sum_{m_f > m_H} Q_f^2\right)^2 \left(\frac{m_H}{10\,\text{GeV}}\right)^3$
1.1 MeV – 0.2 GeV	$H^0 \to e^+e^-$	$\frac{G_F m_e^2}{4\sqrt{2}\pi}\left(1 - \frac{4m_e^2}{m_H^2}\right)^{3/2} \approx 1.7 \cdot 10^{-5}\,\text{eV} \cdot \left(\frac{m_H}{100\,\text{MeV}}\right)$ $\tau_{H^0 \to e^+e^-} \approx 3.8 \cdot 10^{-11}\,\text{s} \cdot \left(\frac{100\,\text{MeV}}{m_H}\right)$
0.21–1 GeV	$H^0 \to \mu^+\mu^-$	$\frac{G_F m_\mu^2}{4\sqrt{2}\pi} m_H \left(1 - \frac{4m_\mu^2}{m_H^2}\right)^{3/2} \approx 7\,\text{eV} \cdot \left(\frac{m_H}{1\,\text{GeV}}\right)$
	$H^0 \to \pi^+\pi^-, \pi^0\pi^0$	$\Gamma(H^0 \to \pi^+\pi^-) = 2\Gamma(H^0 \to \pi^0\pi^0) \sim 0.1-1$ eV (see (1.16))
1–4 GeV	$H^0 \to s\bar{s} \to K\bar{K}$, $K^*\bar{K}^*$, ...	$40\,\text{eV} \cdot \left(\frac{m_H}{1\,\text{GeV}}\right)$
4–10 GeV	$H^0 \to \tau^+\tau^-$	$20\,\text{keV} \cdot \left(\frac{m_H}{10\,\text{GeV}}\right)$
	$H^0 \to c\bar{c}$	$40\,\text{keV} \cdot \left(\frac{m_H}{10\,\text{GeV}}\right)$
10 GeV–$2m_t$	$H^0 \to b\bar{b}$	$0.45\,\text{MeV} \cdot \left(\frac{m_H}{10\,\text{GeV}}\right)$
$>2m_t$	$H^0 \to t\bar{t}$	$290\,\text{MeV} \cdot \left(\frac{m_t}{40\,\text{GeV}}\right)^2 \left(\frac{m_H}{100\,\text{GeV}}\right)$
$>2m_W, m_Z$	$H^0 \to W^+W^-, Z^0Z^0$	$\Gamma_{W^+W^-, Z^0Z^0} > 1.5$ GeV for $m_H > 200$ GeV, $\Gamma_{W^+W^-, Z^0Z^0} \approx 61\,\text{GeV} \cdot \left(\frac{m_H}{500\,\text{GeV}}\right)^3$ for $m_H \gg 2m_Z$ (see (1.17))

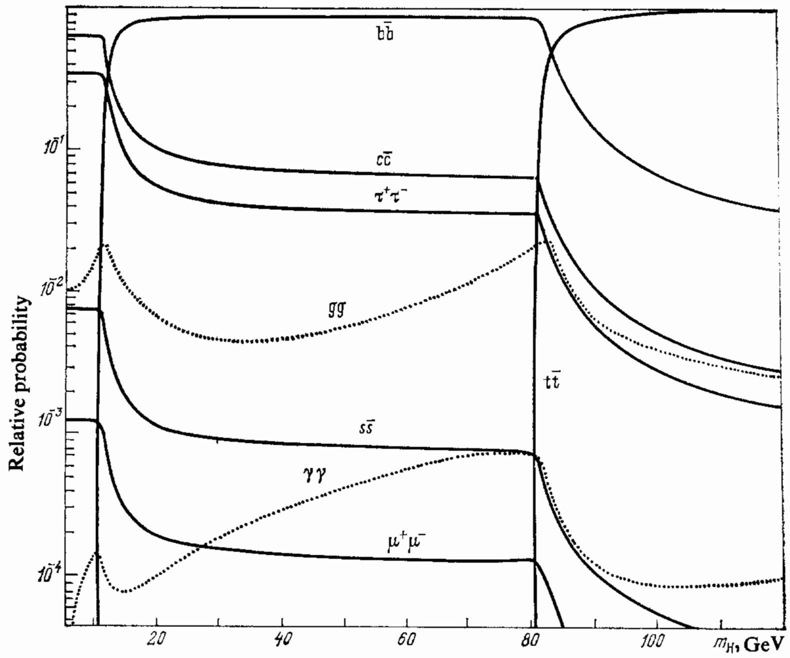

FIG. 4. Relative probabilities for H^0 decays for various values of the mass m_H. The value $m_t = 40$ GeV is assumed.

$$\Gamma(H^0 \to 2\gamma) \approx 0.11\,\mathrm{eV}\cdot\Bigl(7-\frac{4}{3}\sum_{m_f>m_H} Q_f^2\Bigr)\Bigl(\frac{m_H}{10\,\mathrm{GeV}}\Bigr)^3,$$

$$\Gamma(H^0 \to 2g) \approx 0.17\,\mathrm{keV}\cdot\Bigl(\frac{\alpha_s(m_H^2)}{0.15}\Bigr)^2 n_h^2\Bigl(\frac{m_H}{10\,\mathrm{GeV}}\Bigr)^3. \tag{1.18}$$

The widths of the basic decay modes are summarized in Table I; the relative probabilities for the decays are shown in Fig. 4 (the mass of the t quark is assumed to be[2)] 40 GeV).

2. PRODUCTION OF THE STANDARD HIGGS BOSON

Many experimentalists and theoreticians have now taken up the search for Higgs particles. However, it is not by chance that these particles have been called "elusive," Their production cross sections are generally quite small, and to identify them is an extremely complicated experimental problem, which is generally, aggravated by a substantial background. This elusiveness of the Higgs bosons from the experimental standpoint can be explained by their extremely weak coupling with ordinary quarks and leptons and by the particular way in which they decay, giving rise to many-particle final states. It should also be emphasized that both the cross sections for the production of H^0 bosons and the identification methods depend strongly on the mass of these bosons.

In this section we shall examine the reactions which are presently regarded as the most promising from the standpoint of the production of neutral Higgs bosons of the minimal Glashow-Weinberg-Salam scheme.

a) Associative producion of H⁰ with gauge W and Z bosons

It was recognized a comparatively long time ago that the "bremsstrahlung" of H^0 particles by gauge bosons has many advantages for an effort to observe neutral Higgs particles (Fig. 5). Corresponding to this bremsstrahlung are large three-boson vertices:

$$g_{HZZ} = 2\sqrt{G_F\sqrt{2}}\,m_Z^2,\quad g_{HWW} = 2\sqrt{G_F\sqrt{2}}\,m_W^2.$$

Advantages of this mechanism for H^0 production are the comparatively large values expected for the cross sections and the convenient identification conditions. For this reason, it is the associative productions of H^0 with gauge bosons that is regarded as the best bet for an early discovery of Higgs particles in e^+e^- collisions (Fig. 5a) and hadron collisions (Fig. 5b).

Let us examine the characteristic properties of the associative production of H^0.

1) Production of H⁰ *in decays of the* Z⁰ *boson*[35,37]

We consider the decays

$$Z^0 \to H^0 + Z^0_{\text{virt}},\ \text{где}\ \ f\bar{f} = q\bar{q},\ l^-l^+,\ \nu\bar{\nu}\ (l^\pm\ —$$
$$\quad\ \ \llcorner\!\!\longrightarrow t\bar{t}$$

where $f\bar{f} = q\bar{q}, l^-l^+, \nu\bar{\nu}$ ($l^\pm$ is a charged lepton). The propagator of the virtual Z^0 gives rise to a sharp peak in the energy distribution of the Higgs boson at a low energy $E_H = x_H m_{Z/2}$ ($2\mu_H \leqslant x_H \leqslant 1+\mu_H^2$, $\mu_H = m_H/m_Z$). The fermions f and $\bar{f}$ are emitted predominantly in opposite directions, and their binary mass $m_{f\bar{f}}$ tends toward a maximum

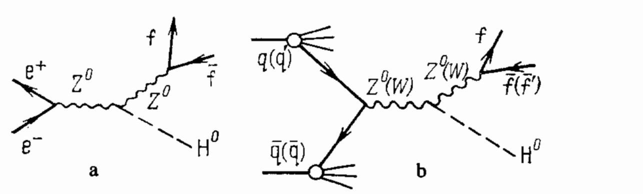

FIG. 5. Feynman diagrams of the associative production of H^0 with Z^0 and W bosons in (a) e^+e^- annihilation and (b) a hadron-hadron collision.

[2)]According to preliminary reports,[137] the UA1 group has discovered the t quark, with a mass of 40 ± 10 GeV on the Sp$\bar{p}$S collider.

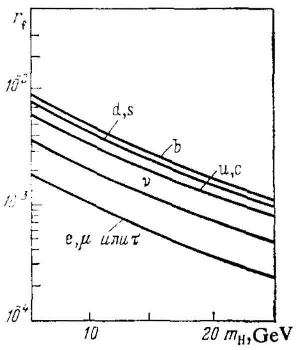

FIG. 6. The ratios $r_f = \dfrac{[\Gamma(Z^0 \to H^0 + f\bar{f})]}{[\Gamma(Z^0 \to \mu^+\mu^-)]}$ as functions of the mass m_H at $\sin^2\theta_W = 0.23$.

value $m_Z - m_H$. This decay could be identified conveniently by (for example) detecting the f and $\bar{f}$ and reconstructing the missing mass m_H. The distribution in x_H in the case of the lepton decay $Z^0 \to H^0 + e^+e^-$ $(\mu^+\mu^-)$ is[37]

$$\frac{1}{\Gamma(Z^0 \to e^+e^-)} \frac{d\Gamma(Z^0 \to H^0 + e^+e^-)}{dx_H} = \frac{\alpha}{\pi} \frac{1}{\sin^2 2\theta_W} \rho(x_H) \frac{\sqrt{x_H^2 - 4\mu_H^2}}{(x_H - \mu_H^2)^2}, \qquad (2.1)$$

where

$$\rho(x_H) = 1 - x_H + \frac{x_H^2}{12} + \frac{2}{3}\mu_H^2, \qquad (2.2)$$

and $\Gamma(Z^0 \to e^+e^-) \approx 90$ MeV is the width of the lepton decay of Z^0. The distributions in $m_{e^+e^-}$ in the lepton cascade are analyzed in detail in Ref. 38.

Figure 6 shows the dependence on the mass of the H^0 boson of the width ratio

$$r_l = \frac{\Gamma(Z^0 \to H^0 + l^+ + l^-)}{\Gamma(Z^0 \to \mu^+\mu^-)}$$

for $m_H \leqslant 25$ GeV. Also shown here are the ratios r_f for other cascades:

$$Z^0 \to H^0 + Z^0_{\rm virt}\,. \\ \quad\quad\quad\quad\ \ \llcorner\!\to f\bar{f}$$

If we assume the standard value $B(Z^0 \to e^+e^-) \approx 3\%$, with $m_H = 10$ GeV we would have $B(Z^0 \to H^0 + e^+e^-) \approx 3 \cdot 10^{-5}$. At larger values of m_H, this quantity decreases rapidly,

$$B(Z^0 \to H^0 + e^+e^-) \approx 10^{-6}$$

at $m_H \approx m_Z/2 \approx 50$ GeV.

The total relative probability for H^0 production in the decays of the Z^0 boson,

$$B(Z^0 \to H^0 + \text{all}) = \sum_f B(Z^0 \to H^0 + f\bar{f}),$$

varies from $\sim 10^{-3}$ to 10^{-4} in the mass interval $m_H = 10$–30 GeV.

The relative probability for a cascade neutrino transition

$$B(Z^0 \to H^0 + \sum_e \nu_e\bar{\nu}_e) \approx 6B(Z^0 \to H^0 + e^+e^-). \qquad (2.3)$$

There is the hope that this transition might be detected in calorimetric measurements.

An important property of the decay $Z^0 \to H^0 + Z^0_{\rm virt}$ is the relatively large fraction of events involving the production of four heavy entities (Q quarks or τ leptons) in the final state.[35,36] The reason for this property is the fact that (if the transition $Z^0_{\rm virt} \to t\bar{t}$ is ignored) the $Z^0_{\rm virt}$ will undergo a transition to $c\bar{c}$, $b\bar{b}$, or $\tau^+\tau^-$ in about a third of the cases. The comparatively long lifetimes of the heavy particles $(\sim 3 \cdot 10^{-13}$–10^{-12} s) and the paticular nature of the fragmentation of the heavy Q quarks—which has the consequence that most of their momentum (as calculated in quantum chromodynamics; see Ref. 40, for example) is carried off by heavy hadrons—make a search for such cascades extremely attractive. The use of a Z^0 "factory" to search for Higgs particles appears promising up to $m_H \lesssim 50$ GeV.

2) $e^+e^- \to Z^0_{\rm virt} \to H^0 + Z^0$

At e^+e^- collision energies $\sqrt{s} > m_Z + m_H$, the binary production of real Z^0 and H^0 particles through a virtual Z^0 boson is extremely promising[8,35,36] (Fig. 5a). The total cross section for this process is

$$\sigma(e^+e^- \to H^0 + Z^0) = \frac{G_F^2 m_Z^4}{96\pi}\left(\frac{1+v_e^2}{s}\right) f(s,\ m_Z^2,\ m_H^2), \qquad (2.4)$$

where

$$f(s,\ m_Z^2,\ m_H^2) = \varkappa(\varkappa^2 + 12y)\frac{1}{(1-y)^2}; \qquad (2.5)$$

here $y = m_Z^2/s$, $\varkappa = 2k/\sqrt{s}$ (k is the c.m. momentum of the H^0 boson), and $v_e = 4\sin^2\theta_W - 1$.

An important property of this process is the fact that at $m_H < m_Z$ the cross section for this process is comparable to

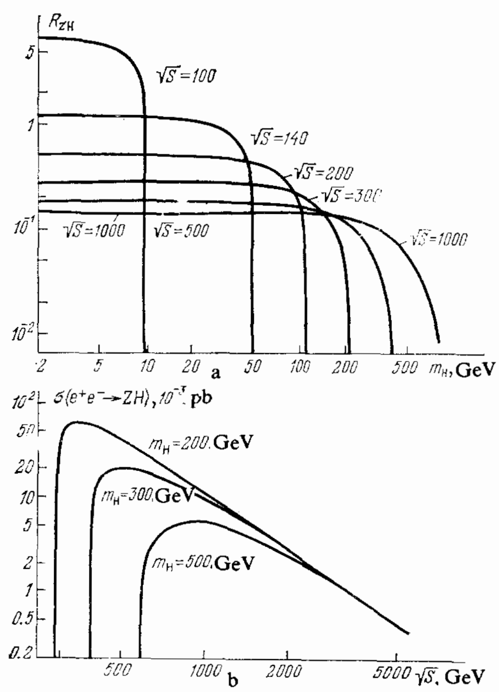

FIG. 7. The ratio R_{ZH} (a) and the cross section $\sigma(e^+e^- \to H^0 + Z^0)$ (b) at various energies as functions of the mass of the H^0 boson for $\sin^2\theta_W = 0.23$.

or even greater than the standard electromagnetic cross section for the process $e^+e^- \to \gamma_{virt} \to \mu^+\mu^-$,

$$\sigma_{pt} = \frac{4\pi\alpha^2}{3s} = \frac{87}{s\,(\mathrm{GeV}^2)}(\mathrm{nb}); \tag{2.6}$$

and is essentially independent of the mass m_H nearly all the way to the production threshold (Fig. 7). If we write the ratio of cross sections (2.4) and (2.6) as

$$R_{ZH} = \sigma(e^+e^- \to H^0Z^0)\,\sigma_{pt}^{-1} = 0{,}125 f(s,\ m_Z^2,\ m_H^2), \tag{2.7}$$

then we find that at $s \gg s_{threshold} = (m_Z + m_H)^2$ we have $f \to 1$. At $m_H \lesssim 2m_Z$ the function f has a maximum at a certain s_m, and the condition $f(s_m) > 1$ holds. At massas $m_H \lesssim 30$ GeV, this maximum becomes extremely sharp and is reached at $\sqrt{s_m} \approx m_Z + \sqrt{2}m_H$. Here we have $f(s_m) \approx 3m_Z/m_H$ and $R_{ZH} \approx 34/m_H$ (GeV). In the case $m_H > 2m_Z$ the function f increases monotonically beginning at $s = s_{threshold}$, and at $m_H \gg m_Z$ we have $f \approx (2k/\sqrt{s})^3$. The cross section $\sigma(e^+e^- \to H^0Z^0)$ itself has a very broad maximum in this case, at about $\sqrt{s} \approx \sqrt{s'_m} \approx 2m_H - m_Z$.

The production of the system $Z^0 + H^0$ can be clearly identified experimentally by making use of the characteristic decay modes of Z^0 and/or H^0 (Ref. 35). In a study of, say, the reaction

$$e^+e^- + \underbrace{Z^0}_{\longrightarrow \mu^+\mu^-} \qquad + \underbrace{H^0}_{\longrightarrow \text{all}} \tag{2.8}$$

the experiment should reveal a peak in the missing-mass spectrum corresponding to the $\mu^+\mu^-$ pair.

Let us compare the characteristics of the processes

$$e^+e^- \to Z^0 \to \underbrace{Z^0_{virt}}_{\longrightarrow \mu^+\mu^-} + H^0,$$

$$e^+e^- \to Z^0_{virt} \to \underbrace{Z^0}_{\longrightarrow \mu^+\mu^-} + H^0.$$

Figure 8 shows the energy dependence of the total cross section for the process $e^+e^- \to H^0 + \mu^+\mu^-$ for the case

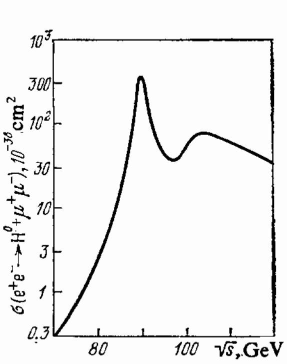

FIG. 8. Energy dependence of the total cross section for the production $e^+e^- \to H^0 + Z^0 \to H^0 + \mu^+\mu^-$ with $m_H = 10$ GeV and $\sin^2\theta_W = 0.23$ (Ref. 43).

$m_H = 10$ GeV at values $\sqrt{s} \gg m_Z$, i.e., beginning in the region of the Z^0 resonance discussed above (modifications stemming from radiation effects in the production of the Z^0 boson are ignored here; see Ref. 41 and 42, for example). The second maximum in this figure is near $\sqrt{s} \approx m_Z + \sqrt{2}m_H$ when the ratio R_{ZH} is at a maximum.

It is not difficult to see that the cross section for $H^0\mu^+\mu^-$ production beyond the threshold for H^0Z^0 production is about an order of magnitude smaller than the resonant cross section. The maximum in the distribution in the invariant mass of the lepton pair, $m_{l^+l^-}$, is, however, much more sharply defined in the case $e^+e^- \to Z^0_{virt} \to Z^0H^0$ (Ref. 43). This circumstance substantially improves the background situation in observations of events of the type in (2.8).

The process $e^+e^- \to H^0Z^0$ is thus an extremely promising field for a search for Higgs bosons in the comparatively broad mass interval $m_H \lesssim \sqrt{s}/2$. A study of the region of larger H^0 masses, $m_H \gg m_Z$, would require that the luminosity of the accelerator permit reliable measurements of cross sections at the level $\sim 0.01\ \sigma_{pt}|_{\sqrt{s} \sim m_H}$.

3) $p + \bar{p} \to H^0 + W^{\pm}(Z^0) + \ldots$

The associative mechanism for H^0 production in hadron collisions is related to elementary quark-annihilation processes

$$q'(q) + \bar{q} \to W^{\pm}(Z^0) \to H^0 + W^{\pm}(Z^0) \tag{2.9}$$

(Fig. 5b), which are described by formulas similar to those discussed in the preceding section in an analysis of the process $e^+e^- \to H^0 + Z^0$. The corresponding cross sections can be calculated[36] in the spirit of the classical Drell-Yang model by introducing distribution functions of the u and d quarks in nucleons. In particular, the following result has been derived:

$$\sigma_{DY}(p\bar{p} \to H^0 + W^-(W^+) + \text{all}) = \frac{G_F^2 m_W^4}{72\pi s} \int_{\tau_0}^{1} \frac{dx_1}{x_1} \int_{\tau_0/x_1}^{1} \frac{dx_2}{x_2}\, u(x_1)\, d(x_2)\, f(x_1x_2s,\ m_W^2,\ m_H^2), \tag{2.10}$$

where $\tau_0 = (m_W + m_H)^2/s$, and the function $f(s, m^2, m_H^2)$ is defined in (2.5). The distributions of the u and d quarks in the proton are normalized by

$$\int_0^1 (u(x) - \bar{u}(x))\,dx = 2, \qquad \int_0^1 (d(x) - \bar{d}(x))\,dx = 1.$$

If m_H is not too large, the quantum-chromodynamics effects which modify the Drell-Yang formulas become significantly weaker in a calculation of the ratio of the $H^0W^{\pm}(Z^0)$ and $W^{\pm}(Z^0)$ yields. This comment applies primarily to the multiplicative K factor known from an analysis of lepton-pair production.

Figure 9 shows results (taken from Ref. 36) calculated for the ratios

$$\rho_H^W = \frac{\sigma(p\bar{p} \to H^0 + W^{\pm} + \ldots)}{\sigma(p\bar{p} \to W^{\pm} + \ldots)}, \quad \rho_H^Z = \frac{\sigma(p\bar{p} \to H^0 + Z^0 + \ldots)}{\sigma(p\bar{p} \to Z^0 + \ldots)}$$

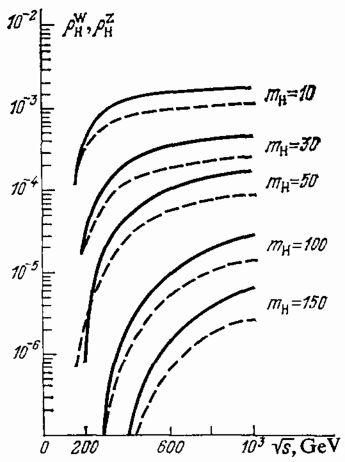

FIG. 9. Energy dependence of the ratios ρ_H^W (solid curves) and ρ_H^Z (dashed) for $\sin^2\theta_W = 0.25$ (Ref. 36).

of the cross sections for $H^0W^\pm$ (Z^0) production to the cross sections for the single production $W^\pm$ (Z^0) in $p\bar{p}$ collisions. It follows from this figure that the cross section for associative production falls off rapidly with increasing mass. This figure also clearly demonstrates the threshold effects at various values of m_H. For m_H between 10 and 50 GeV, in the energy range of the Sp$\bar{p}$S collider at CERN, the ratios $\rho_H^{W,Z}$ are, in order of magnitude, B ($Z^0 \to H^0 + \ldots$) $\sim$ B ($W^- \to H^0 + \ldots$) $\sim 10^{-3} - 10^{-4}$. The cross sections for H^0W^- and H^0Z^0 production here would be $\sim$3–0.2 pb and $\sim$0.8–0.05 pb. The energies required to reach the same cross sections in pp scattering would be greater: $\sqrt{s} \gtrsim 800$ GeV.

It is not difficult to see that the cross sections for $p\bar{p} \to H^0Z^0$ + all at the energies of the Sp$\bar{p}$S collider would be about 1.5–2 orders of magnitude smaller than the corresponding cross sections for the production of H^0 in e^+e^- annihilation near the threshold. The cross sections for H^0Z^0 production in $p\bar{p}$ (pp) and e^+e^- collisions should become comparable only when the energy of the hadronic process, $\sqrt{s}$, exceeds $\sqrt{s_m}$ by more than an order of magnitude (we recall that the average fraction of the proton energy per valence quark is $\langle x_q \rangle \approx 0.1$). When effects of the breaking of scaling in the structure functions are taken into account, and these effects are particularly important at large masses m_H, we find some further reductions in the cross sections in hadron collisions.

Since the cross sections are so small, it is unlikely that the cascades $Z^0 \to H^0 + l^+l^-$ in $\bar{p}$(p)p collisions will be observed in the near future. For the process $\bar{p}p \to W^\pm + H^0$ + all, with a cross section about an order of magnitude larger than that for H^0Z^0, the situation is more favorable. This cross section is comparable to $\sigma(e^+e^- \to H^0Z^0)|_{s=s_m}$ at $\sqrt{s} \approx 4\sqrt{s_m}$. Hopes for discovering standard H^0 bosons on the Sp$\bar{p}$S collider are pinned on the cascades

$$W \to H^0 + W \to H^0 + f'\bar{f}\,(q'\bar{q},\ l\bar{\nu}_e), \qquad (2.11)$$

with H^0 being identified on the basis of the heavy particles, and $W^\pm$ being identified on the basis of their decay products (lepton and quark jets).

Let us examine some general properties of cascade (2.11). In the distribution in the invariant mass of the $f'\bar{f}$ pair, $m_{f'f}$, there are two rather well-defined kinematic regions (Fig. 10; cf. Fig. 8). One (the humps at the left in Fig. 10) corresponds to the transition

$$\begin{array}{l} W \to W_{\rm virt} \quad + H^0. \\ \quad\ \ \lfloor\!\longrightarrow f'\bar{f} \end{array}$$

In this region, as was mentioned earlier, the propagator of the virtual boson has the consequence that the mass $m_{ff'}$ tends toward its kinematic limit: $m_{ff'} \to m_W - m_H$. The other region is associated with the cascade

$$\begin{array}{l} W_{\rm virt} \to W \quad + H^0, \\ \qquad\quad \lfloor\!\longrightarrow f'\bar{f} \end{array}$$

and this is the region of the sharp peaks in Fig. 10, which are related to the propagator of the external W.

The expected values of the cross sections for $W^\pm H^0$ production, with allowance for the proposed modifications of the Sp$\bar{p}$S collider and the experimental apparatus (Ref. 44,

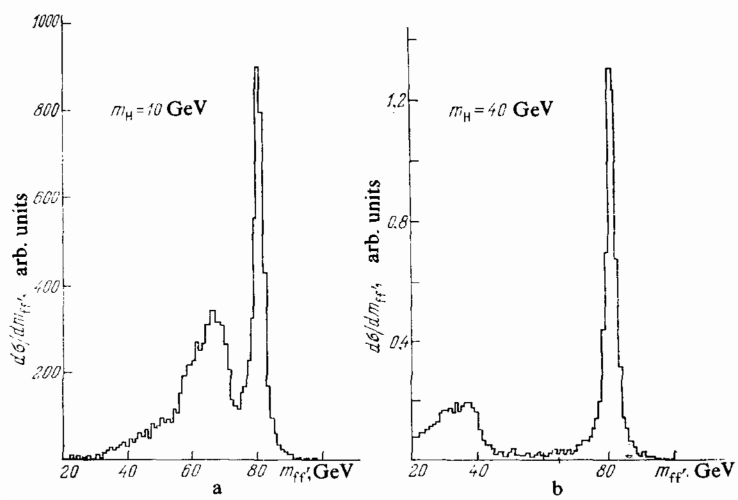

FIG. 10. Expected distributions in the invariant mass $m_{ff'}$ in the process $p\bar{p} \to W_{\to ff'} + H^0$ + all at $\sqrt{s}$ 660 GeV (Ref. 44). a—$m_H = 10$ GeV; b—40 GeV.

for example), raise the hope that the first searches for associative production of Higgs bosons with masses of the order of tens of GeV can be carried out on the Sp$\bar{p}$S collider. Furthermore, if an H^0 boson with a mass $m_H \sim 10$ GeV does exist, then it is quite possible that it has already been detected in experiments on the Sp$\bar{p}$S collider, but it has not yet been extracted from the data and thus has not been recognized. The statistical base of ~ 100 events of $W^+ \to e^\pm \nu_e$ decays[1,2] which has been built up corresponds to the production of $\sim 10^3 W^\pm$ particles and therefore roughly one event involving the production of a $W^\pm H^0$ state.

The hope for studying heavy H^0 bosons in $\bar{p}p$ scattering is pinned on a substantial increase in the energy of the colliding beams. For example, for $m_H = 200$, 300, and 400 GeV the cross sections for H^0W^- production, even with $\sqrt{s} = 2$ TeV, are ~ 0.15 pb, $2.5 \cdot 10^{-2}$ pb, and $6 \cdot 10^{-3}$ pb. At $m_H = 500$ GeV and $\sqrt{s} = 5$ TeV, the cross section is $\sim 4.4 \cdot 10^{-2}$ pb. It follows from (2.10) and (2.5) that at masses m_H far greater than m_W the cross section for associative production should be described approximately by the scaling formula

$$\sigma_{HW}(m_H, \sqrt{s}) = \frac{m_0^2}{m_H^2}\sigma_{HW}\left(m_0, \sqrt{s}\frac{m_0}{m_H}\right), \qquad (2.12)$$

where $\sigma_{HW}(m_0, \ldots)$ is the cross section for the production of the W^-H^0 system in $\bar{p}p$ collisions at a certain fixed mass of the Higgs boson, $m_0 \gg m_W$. The estimates generated by this formula can be used to scale up the known cross sections to higher energies $\sqrt{s}$ and larger masses H^0.

If $m_H > 2m_W$, this mechanism will lead to an extremely distinctive signal, corresponding to the production of three W bosons or $2Z^0$ and W in the final state.[3)]

b) "Gluon" production of H^0 in $p\bar{p}$ and pp collisions

There may be a direct production of H^0 in high-energy hadron-hadron collisions through the annihilation of a pair of gluons into a Higgs boson (Fig. 11). The cross section for H^0 production by the gluon mechanism turns out to be comparatively large. The reason is that, although the amplitude for $H^0 \to 2g$ explicitly contains the small factor $\alpha_s(m_H^2)/6\pi$, it is (roughly speaking) proportional to the invariant mass of the two gluons, according to (1.5a). The cross section for the reaction pp ($\bar{p}$)$\to$hadrons is[45]

$$\sigma(pp \to H^0 + \ldots) = \sigma_0^H n_h^2 L, \qquad (2.13)$$

where

$$\sigma_0^H = \frac{G_F}{\sqrt{2}}\frac{\pi}{288}\left(\frac{\alpha_s(m_H)}{\pi}\right)^2 \approx \left(\frac{\alpha_s(m_H)}{0{,}15}\right)^2 \cdot 8 \cdot 10^{-38}\ \text{cm}^2. \qquad (2.14)$$

The quantity n_h is actually the number of heavy quarks with masses $m_Q > 0.2\, m_H$ (Sec. 1). In the parton model the function L is

[3)]M. B. Voloshin and L. V. Okun' have also discussed the possibility of observing the associative production of H^0 bosons on the basis of the inclusive production of three gauge bosons.

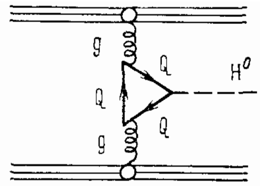

FIG. 11. Feynman diagram corresponding to a gluon mechanism for the production of an H^0 boson in a hadron-hadron collision.

$$L = L(\tau) = \tau \int dx_1 dx_2 g(x_1)\, g(x_2)\, \delta(x_1 x_2 - \tau); \qquad (2.15)$$

here $\tau = m_H^2/s$, and $g(x)$ is the distribution function of the gluons in a proton ($\int_0^1 g(x)\, x\, dx \approx 0.5$). In a realistic case, expression (2.15) would have to be modified for quantum-chromodynamics effects, in particular, the effects which stem from the known deviations from scaling in the functions $g(x)$ and from preasymptotic effects.

Figure 12 shows estimates of the cross section for H^0 production in this reaction as a function of m_H for $\sqrt{s} = 400$ GeV and $\sqrt{s} = 2000$ GeV according to Ref. 46, where the assumption $n_h = 3$ was used. At $\sqrt{s} \approx 540$ GeV and $m_H = 10$ GeV the cross section is ~ 40 pb. This value is about an order of magnitude greater than the cross section for the associative production of an H^0 boson with W^-, and it turns out to be comparable to the cross section

$$\sigma(\bar{p}p \to \underbrace{Z^0}_{\to e^+e^-} + \ldots) \approx 30\ \text{pb},$$

which is presently being measured successfully (admittedly, under good background conditions). Consequently, in the experiments on the Sp$\bar{p}$S collider, where ten $Z^0 \to e^+e^-$ events have been observed, rather light Higgs bosons with $m_H \lesssim 20$–30 GeV should already have been produced. Unfortunately, at $m_H < 2m_W$ in the gluon production of H^0 we do not have the convenient way to identify the Higgs boson that we have in the case of its associative production with $W^\pm$ and Z^0, so that it would be an extremely complicated matter to use gluon production. The one possibility here might be to identify the heavy leptons or quarks into which the H^0 decays, but the background production of $\tau^+\tau^-$ pairs

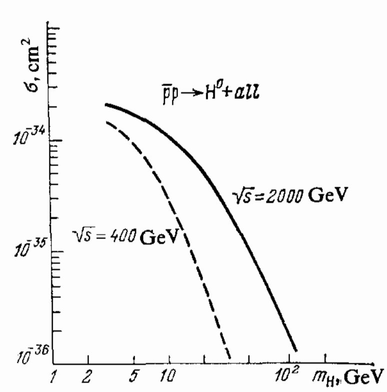

FIG. 12. Total cross section for the process $pp \to H^0 +$ all as a function of the mass m_H at $\sqrt{s} = 400$ and 2000 GeV (Ref. 46).

or heavy quarks in Drell-Yang processes would lead to cross sections considerably larger than that for the production of a Higgs boson. Because of the complexity of identifying the Higgs signal, this reaction is not presently regarded as a good candidate for a search for comparatively light H^0 bosons. For Higgs bosons with $m_H \gtrsim 180$ GeV the gluon mechanism might become a genuinely effective method for searching for Higgs bosons in $\bar{p}p$ and (especially) pp collisions (where the background is lower, but the associative production of H^0 is slight), thanks to the dominant characteristic decay modes $H^0 \to W^+W^-$, Z^0Z^0.

The values expected for the cross sections for the gluon production of H^0 depend strongly on the mass spectrum of heavy quarks, which determines the value of n_h [see Eq. (1.5) in Sec. 1]. Here we will quote some estimates for the case $n_h = 1$. They turn out to be significantly too high if the spectrum of quarks not yet discovered is limited to the t quark with $m_t \ll 0.2\, m_H$, or they turn out to be significantly too low if the spectrum of heavy quarks turns out to be richer.

For $n_h = 1$ the cross sections for the gluon mechanism for H^0 production are about 0.15 pb, $5 \cdot 10^{-2}$ pb, and 2×10^{-2} pb for $\sqrt{s} = 2$ TeV and $m_H = 200, 300$, and 400 GeV. These cross sections are approximately equal to the corresponding cross sections for the associative production of $W^\pm H^0$ (cf. Subsection 2a3). With increasing energy $\sqrt{s}$, the cross sections expected for the gluon production of Higgs bosons begin to exceed σ_{HW} significantly. Expression (2.14) in approximate scaling form is [cf. (2.12)]

$$\sigma_H(m_H, \sqrt{s}) = \left(\frac{\alpha_s(m_H)}{\alpha_s(m_0)}\right)^2 \sigma_H\left(m_0, \sqrt{s}\,\frac{m_0}{m_H}\right), \qquad (2.16)$$

where $\sigma_H(m_0, \sqrt{s})$ are the cross sections for the gluon production of an H^0 boson with a fixed mass m_0 (under the assumption $n_h = \text{const}$). In view of the substantial uncertainties regarding the choice of n_h, the gluon distribution functions, and the quantum-chromodynamics effects, expression (2.16) is completely reasonable for generating preliminary estimates of the cross section σ_H.

If all the quarks which exist in nature have masses $m_Q < 0.2 m_H$, small factors of the type m_Q^4/m_H^4, associated with the form factor at the Hgg vertex,[17,18] will appear in cross sections (2.14) and (2.16) (Sec. 1).

In the case $m_H \lesssim 2m_W$ the decay of H^0 into gauge bosons might be seen in the processes $H^0 \to W^\pm + W^\mp_{virt}$ and $H^0 \to Z^0 + Z^0_{virt}$. Particularly interesting from the standpoint of observing H^0 is the cascade

$$\begin{tikzcd} H^0 \arrow[r] & Z^0 \arrow[d] & + Z^0_{\mathrm{virt}} \arrow[d] \\ & \nu\bar{\nu} & q\bar{q} \end{tikzcd} ,$$

which should lead to the production of hadron jets with a visible imbalance of transverse energy.[4)]

The basic physical backgrounds for the production of heavy bosons with $m_H > 2m_W$ are multijet quantum-chromodynamics processes and reactions of the type $p\bar{p}$ $(p) \to W^+W^-$ (Z^0Z^0) + all. Analysis of the situation (Ref. 47, for example) leaves the hope that at a sufficiently high luminosity and at a sufficiently high energy of the hadron beams the decays of heavy H^0 particles into gauge bosons could be identified comparatively reliably.

In addition to the standard Higgs boson, the gluon mechanism could give rise to other scalar or pseudoscalar particles which are predicted theoretically. For example, the various neutral pseudo-Goldstone particles P^0 which arise in the technicolor theories also have comparatively large P^0gg vertices, of the order in (1.5a). These vertices contain an additional factor N_{TC} (the number of technicolors). These theories predict large cross sections for the production of neutral entities belonging to a colored octet (Ref. 11, for example).

c) The decay $Z^0 \to H^0 + \gamma$

A completely unambiguous identification of a Higgs boson might be achievable in the decay $Z^0 \to H^0 + \gamma$, where a monochromatic photon is emitted. However, there is no tree vertex $Z^0H^0\gamma$, and the process would go only by virtue of W-boson and fermion loops (Fig. 13). The width $\Gamma(Z^0 \to H^0 + \gamma)$ is thus extremely small[49]:

$$\frac{\Gamma(Z^0 \to H^0 + \gamma)}{\Gamma(Z^0 \to e^+e^-)} \approx 8 \cdot 10^{-5}\left(1 - \frac{m_H^2}{m_Z^2}\right)^3 \qquad (2.17)$$

[the contributions of fermion loops (Fig. 13b) have been omitted under the assumption $(m_f/m_Z)^2 \ll 1$]. Measurements of $\Gamma(Z^0 \to H^0 + \gamma)$ would be of interest in connection with attempts to resolve the question of the number of generations of superheavy fermions and also to test the gauge structure of the theory (by virtue of the contribution of W loops). The amplitude for the decay of a Z^0 boson into a γ ray and a neutral technipion would, generally speaking, be of the same order of magnitude as that for $Z^0 \to H^0 + \gamma$, and it would also contain an additional factor of N_{TC}—the number of technicolors.

d) The Higgs boson in the decays of heavy quarkonia

Since the interaction of a Higgs boson with fermions is proportional to the mass of the fermions, it is natural to seek the Higgs boson in the decays of heavy quarkonia. The most promising direction is to search for the H^0 boson in the radiative decays of the vector states $V_Q \equiv {}^3S_1\,(Q\bar{Q})$ (Fig. 14). The probability for such a decay is given by[50]

$$\frac{\Gamma(V_Q \to H^0 + \gamma)}{\Gamma(V_Q \to \mu^+\mu^-)} = \frac{G_F m_V^2}{4\sqrt{2}\pi\alpha}\left(1 - \frac{m_H^2}{m_V^2}\right). \qquad (2.18)$$

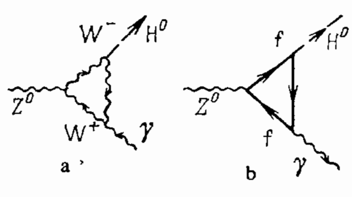

FIG. 13. Feynman diagrams describing the radiative decay $Z^0 \to H^0 + \gamma$ due to a W-boson loop (a) and a fermion loop (b).

[4)]Such cascades were also discussed in Ref. 48 in connection with attempts to explain the "exotic events" observed on the $Sp\bar{p}S$ collider. However, the enhancement by a factor of 10^3 of the Hgg vertex of the hypothetical Higgs boson with a mass of 150 GeV which was proposed in Ref. 48 would, if taken literally, lead to a width $\Gamma(H \to gg)$ greater than m_H.

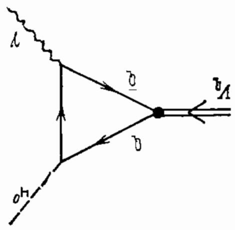

FIG. 14. Radiative decay of vector quarkonium.

Numerically, with $m_H^2/m_V^2 \ll 1$, we find $B(J/\psi \to H^0\gamma) \approx 6.4 \cdot 10^{-5}$ and $B(\Upsilon \to H^0 + \gamma) \approx 2.6 \cdot 10^{-4}$ from (2.18). For toponium, we would have

$$\frac{\Gamma(T \to H^0+\gamma)}{\Gamma(T \to \mu^+\mu^-)} \approx 0.58 \left(\frac{m_T}{80}\,\mathrm{GeV}\right)^2.$$

The total width of toponium, Γ_T, is determined to a large extent by weak interactions, and it depends strongly on the mass m_T. The theoretical calculations of Γ_T are comparatively reliable over a broad range of the mass m_T (Refs. 51 and 52, for example). In particular, at $m_T \gtrsim 45$ GeV we have

$$B(T \to H^0+\gamma) \geqslant 10^{-2}\left(1-\frac{m_H^2}{m_T^2}\right).$$

From the standpoint of a search for a Higgs boson in radiative decays of toponium, the experiments of greatest interest today are those involving e^+e^- beams directly at the T resonance. In the $\bar{p}p$ and pp collisions, the radiative decays of toponium (with $m_T < m_Z$), which constitute a significant fraction of the traditional $\mu^+\mu^-$ signal, whould give rise to a solitary photon with a large $p_\perp$. The corresponding cross sections for the production $H^0 + \gamma$ at energies $\sqrt{s}$ in the Sp$\bar{p}$S-collider range are expected to be of the order of a fraction of a picobarn. The background situation, however, involving large contributions from direct photons, appears to be quite complicated.

In the decays of quarkonia, we might note, it would be possible to observe not only the standard Higgs boson, H^0, but also pseudoscalar particles which are coupled relatively strongly with heavy fermions. If the pseudoscalar P^0 has an interaction of the type $(m_q/v)(\overline{q}i\gamma_5 q)P^0$, with quarks [cf. (1.4)], then both the width $\Gamma(V_Q \to P^0 + \gamma)$ and the angular distribution of the γ ray would be identical for the scalar H^0 and the pseudoscalar P^0. It would be possible to distinguish H^0 from P^0 by, for example, measuring the angular distributions of the lepton pair in the case of conversion of a γ ray.

We will complete this section of the paper with a few comments regarding other mechanisms for the production of H^0 bosons—mechanisms which are not presently regarded as leading candidates. Some of these mechanisms have been discussed elsewhere.[12–14] A Higgs boson might be produced, for example, along with a pair of heavy quarks or leptons through a bremsstrahlung process (Ref. 53, for example). This mechanism, however, would be of practical interest only for the t quark or for even heavier fermions.

We have discussed the process $e^+e^- \to Z^0 \to H^0 + Z^0$ under the assumption that one of the Z^0 bosons is on the mass shell (Subsections 2a1 and 2a2). If the mass of a Higgs boson turned out to be greater than $\sqrt{s} - m_Z$, then the range of m_H values to be studied could be expanded slightly by the process $e^+e^- \to H^0 + f\bar{f}$, where $m_{f\bar{f}} < m_Z$. In addition to the transition

$$e^+e^- \to Z^0_{\mathrm{virt}} \to H^0 + Z^0_{\mathrm{virt}}, \quad Z^0_{\mathrm{virt}} \to f\bar{f} \tag{2.19}$$

an important role might be played here by the production of H^0 in a "collision" of gauge W and Z^0 bosons[54] (Fig. 15). This mechanism could be identified experimentally very well. For the collision of W bosons, the cross section is considerably larger than that for Z^0, proportional to the quantity

$$\sigma_{WW} = \left(\frac{\alpha}{\sin^2\theta_W}\right)^2 \frac{\Gamma(H \to W^+W^-)}{m_H^3} = \frac{\alpha^2 G_F}{8\pi\sqrt{2}\sin^4\theta_W}$$

[cf. (2.14)]. In contrast with transition (2.19), the cross section does not diminish with increasing s far from the threshold. At high energies, $\sqrt{s} \gtrsim 1$ TeV, the reaction $e^+e^- \to \nu\nu H^0$, involving a W-boson mechanism, might become important for the production of H^0 with $m_H \gtrsim 0.4$ TeV. The corresponding cross sections would be $\gtrsim 0.05$ pb at $\sqrt{s} \gtrsim 1$ TeV.

At very large masses m_H, the production of H^0 by a pair of virtual W (Z) bosons in pp and $\bar{p}p$ collisions (Fig. 15c; see, for example, Refs. 47 and 55) might turn out to compete with the gluon mechanism if, for example, there are no heavy quarks in nature with masses $m_Q \gtrsim 0.2m_H$. In particular, if the heaviest quark is the t quark with a mass $m_t \approx 40$ GeV, then the decrease in the form factor would cause the contributions of the gluon and W-boson mechanisms to become comparable at $m_H \approx 400$–500 GeV. It should be kept in mind, however, that these cross sections become equal in approximately that region of masses m_H where the Higgs boson is extremely wide, and its direct observation seems quite problematical.

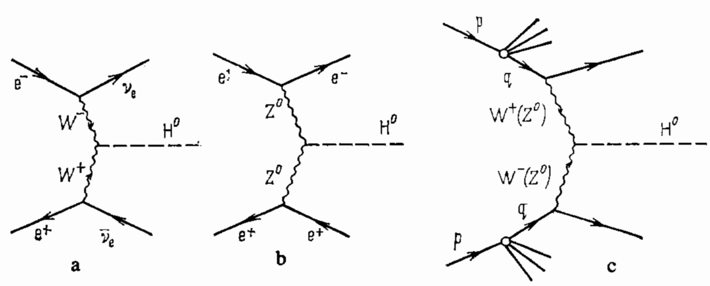

FIG. 15. Production of an H^0 boson by a pair of gauge bosons in (a, b) e^+e^- collisions and (c) pp scattering.

3. NONSTANDARD HIGGS BOSONS

Up to this point we have been discussing the simplest possibility: that the theory contains a single doublet of Higgs fields or, equivalently, that there exists only one physically observable neutral Higgs boson. Actually, it is completely possible that the number of Higgs bosons is not this small. We turn now to certain versions of the theory which require expansions of the Higgs sector.

a) Charged Higgs bosons

We begin with the simplest expansion of the sector of scalar particles: a standard model with several Higgs doublets. On the one hand, this modification does not represent a substantial complication of the minimal standard scheme, and on the other hand doublet Higgs fields by themselves do not alter the relation $\rho = M_W^2/m_Z^2 \cos^2\theta_W = 1$, which stands up well experimentally. The charged Higgs bosons which appear in this model are of interest from the experimental standpoint.

We assume that there are n doublets

$$\varphi_i = \begin{pmatrix} \varphi_i^{(+)} \\ \frac{1}{\sqrt{2}}(v_i + H_i^0 + iP_i^0) \end{pmatrix}, \quad \langle H_i^0\rangle = \langle P_i^0\rangle = 0 \tag{3.1}$$

(we assume that the vacuum expectation values v_i are real). The combinations g^+ and g^0,

$$g^+ = \frac{1}{v}\sum_{i=1}^{n} v_i\varphi_i^{(+)}, \quad g^0 = \frac{1}{v}\sum_{i=1}^{n} v_i P_i^0,$$
$$v^2 = \sum_{i=1}^{n} v_i^2 = (G_F\sqrt{2})^{-1}, \tag{3.2}$$

are nonphysical Goldstone bosons, while the $(n-1)$ charged and $(2n-1)$ neutral particles which are orthogonal to g^+ and g^0 remain in the spectrum of physical states. The interactions of these Higgs bosons are not rigidly fixed, in contrast with the minimal scheme with a single Higgs boson; this comment applies in particular to the interactions of these bosons with fermions. On the other hand, some restrictions are imposed on the form of the Yukawa couplings by the condition that flavors must be "naturally" conserved in an interaction with neutral Higgs particles, i.e., the condition that there are no vertices $H_i^0(\bar{s}d)$, $P_i^0(\bar{s}d)$, $H_i^0(\bar{c}u)$, etc., whose existence would lead to a simulation of unacceptably strong neutral currents with changes in strangeness, charm, etc. This condition requires that no more than three doublets[56] (φ_1, φ_2, and φ_3, say) initially interact with fermions, and the most general form of their interaction would be

$$\begin{aligned} -L = &-\frac{\varphi_1^{(+)}}{v_1}(m_d\bar{u}'_L d_R + m_s\bar{c}'_L s_R + m_b\bar{t}'_L b_R) \\ &+\frac{\varphi_2^{(+)}}{v_2}(m_u\bar{u}_R d'_L + m_c\bar{c}_R s'_L + m_t\bar{t}_R b'_L) \\ &-\frac{\varphi_3^{(+)}}{v_3}(m_e\bar{\nu}_e e_R + m_\mu\bar{\nu}_\mu\mu_R + m_\tau\bar{\nu}_\tau\tau_R) + \text{H.a.} \end{aligned} \tag{3.3}$$

for charged bosons or

$$\begin{aligned} -L = &\frac{1}{v_1}[(m_d\bar{d}d + m_s\bar{s}s + \ldots)H_1^0 \\ &+ (m_d\bar{d}i\gamma_5 d + m_s\bar{s}i\gamma_5 s + \ldots)P_1^0] \\ &+\frac{1}{v_2}[(m_u\bar{u}u + m_c\bar{c}c + \ldots)H_2^0 \\ &-(m_u\bar{u}i\gamma_5 u + m_c\bar{c}i\gamma_5 c + \ldots)P_2^0] \\ &+\frac{1}{v_3}[(m_e\bar{e}e + m_\mu\bar{\mu}\mu + \ldots)H_3^0 \\ &+(m_e\bar{e}i\gamma_5 e + m_\mu\bar{\mu}i\gamma_5\mu + \ldots)P_3^0] \end{aligned} \tag{3.4}$$

for neutral bosons. The vacuum expectation values v_i and the components $\varphi_i^{(+)}$, H_i^0, and P_i^0 are defined in (3.1). In (3.3) we have

$$\begin{pmatrix} d' \\ s' \\ b' \end{pmatrix}_L = V\begin{pmatrix} d \\ s \\ b \end{pmatrix}_L, \quad \begin{pmatrix} u' \\ c' \\ t' \end{pmatrix}_L = V^+\begin{pmatrix} u \\ c \\ t \end{pmatrix}_L, \tag{3.5}$$

where the V_{ij} are the elements of the standard Kobayashi-Maskawa matrix.[57] We should emphasize that any pair of doublets φ_1,φ_2, φ_3—or all three pairs (as in the minimal scheme)—could in fact be a single field. In order to use (3.3) and (3.4) we need to know, in addition to the ratios of the vacuum expectation values v_i, how the physical Higgs bosons with a certain mass are constructed from the fields $\varphi_i^{(+)}$, H_i^0, and P_i^0. Although this "mixing" is determined by the details of the Higgs potential, the expressions given above show that, as for a standard Higgs boson, the constant of the Yukawa interaction of any Higgs particle is proportional to the mass of a fermion (for charged particles, a corresponding element of the Kobayashi-Maskawa matrix also appears).

The structure of interactions (3.3) and (3.4) is actually an extremely general structure. It is reproduced in an arbitrary Higgs sector under the sole requirement of natural conservation of flavors in the exchange of neutral scalar particles. If we ignore a possible large difference between the vacuum expectation values of Higgs fields, and if we also ignore a deviation of the components $\varphi_i^{(+)}$ from the physical charged H^+ states, then the ratios of probabilities for the decays of, say, the H^+ boson by different pathways can be estimated as follows, with an accuracy to phase-volume effects:

$$\Gamma(H^+\to t\bar{b}) : \Gamma(H^+\to c\bar{b}) : \Gamma(H^+\to c\bar{s}) : \Gamma(H^+\to\tau\nu_\tau)$$
$$\approx 1500 : 0.06 : 2 : 1. \tag{3.6}$$

Here we have used the values of the quark mixing parameters[58] and $m_t = 40$ GeV.

It follows from (3.6) that for values of the H^+ mass in the interval 1.8–2.5 GeV the decays of this particle would be determined by the $\tau^+\nu_T$ mode, while at $2m_t > m_{H^+} > 2.5$ GeV the basic decay modes would be $H^+\to\overline{sc}$ and $H^+\to\tau^+\nu_\tau$, which would have some extremely characteristic experimental manifestations. Because of the strong suppression of the charged current $c\bar{b}$ by the matrix element V_{cb}, the decay mode $H^+\to c\bar{b}$ would hardly be important, regardless of the mass of the charged Higgs boson. The H^+ lifetime can be estimated crudely as the lifetime of a standard neutral Higgs boson of similar mass (Table I).

The most obvious manifestation of relatively light charged scalars would be the decay of heavy quarks into these particles and lighter quarks. The amplitude of the corresponding transition is proportional to $m_q/v = \sqrt{G_F\sqrt{2}}\, m_q$; i.e., here we are dealing not with a weak decay but a semiweak decay. If the width of the ordinary decay of, say, the t quark were

$$\Gamma(t\to b q\bar{q},\ bl\nu) \approx \frac{9G_F^2 m_t^5}{192\pi^3},$$

then we would have

$$\Gamma(t\to H^+ + b) \sim \frac{G_F m_t^3}{8\pi\sqrt{2}} \tag{3.7}$$

and the ratio of these widths,

$$\frac{\Gamma(t\to H^+ + b)}{\Gamma(t\to bq\bar{q}',\ bl\nu)} \sim \frac{4\sqrt{2}\pi^2}{3G_F m_t^2}, \tag{3.8}$$

would be of the order of 10^3 if $m_t \sim 40$ GeV. The absolute value of the width would be $\Gamma(t\to H^+ + b) \sim 20$ MeV if $m_t = 40$ GeV or $\Gamma \sim 70$ MeV if $m_t \sim 60$ GeV.

Let us take a brief look at what the existence of charged Higgs bosons $H^\pm$ with a mass $m_{H^+} < m_T/2$ would bode for toponium with a mass $m_T \sim 70$ GeV. In this case the total width of T would be determined by the semiweak decay of a free t quark, $t\to H^+ + b$, within the toponium, and it would reach values of the order of tens of MeV—two or three orders of magnitude larger than the standard expectation[51,52] $\Gamma_T(m_T \sim 70\ \text{GeV}) \sim 50$ keV. At such a large value of Γ_T [which would lead to $B(T\to\mu^+\mu^-) \sim 10^{-3}$] there would be essentially no hope of observing toponium in a hadron-hadron collision on the basis of the standard lepton mode $T\to\mu^+\mu^-$. In the total cross section for e^+e^- annihilation, however, the signal would not actually change since Γ_T would remain below the energy resolution of the e^+ and e^- beams (see Sec. 5 and Table II). Since the decay of H^+ would be dominated by the transition $H^+\to c\bar{s}$, $\tau^+\nu_\tau$, the decay $T\to H^+H^-b\bar{b}$ should be extremely obvious experimentally. In contrast with the standard case (Refs. 41 and 52, for example), the T peak should not be noticeable in the muon mode of e^+e^- annihilation.

The search for charged Higgs bosons is an incomparably simpler experimental task than the search for neutral bosons. For example, if the mass of H^+ turns out to be smaller than m_t, this particle could be easily found by making use of the nonstandard characteristics of the decay of t quarks.

Charged scalar particles might be produced in pairs in e^+e^- collisions. The cross section for their production can be written[59]

$$R^{H^+H^-} = \frac{\sigma(e^+e^-\to H^+H^-)}{\sigma(e^+e^-\to \mu^+\mu^-)} = \frac{\beta_H^3}{4}\left[1 - \frac{2s\varepsilon v_e v_H}{(s/m_Z^2)-1} + \frac{s^2\varepsilon^2(1+v_e^2)v_H^2}{[(s/m_Z^2)-1]^2}\right],$$
$$\beta_H = \sqrt{1-\frac{4m_H^2}{s}},\qquad \varepsilon = \frac{G_F}{8\pi\sqrt{2}\alpha} \approx 4.49\cdot 10^{-5}\ \text{GeV}^{-2},$$
$$v_e = 4\sin^2\theta_W - 1,\quad v_H = -2\cos 2\theta_W. \tag{3.9}$$

At the Z^0 resonance, the ratio of the partial widths of decays by the H^+H^- and $\mu^+\mu^-$ modes is

$$\frac{\Gamma(Z^0\to H^+H^-)}{\Gamma(Z^0\to\mu^+\mu^-)} = \frac{1}{4}\beta_H^3\frac{v_H^2}{1+v_e^2} \approx \frac{1}{4}\beta_H^3. \tag{3.10}$$

According to (3.9), the relative contribution of the binary production of H^+H^- to the total cross section for hadron production in e^+e^- annihilation is numerically small. The best place to observe H^+H^- would be near the threshold, where the overall structure of events should change.

We will also mention the vertices $Z^0H_i^0H_j^0$ and $W^\pm H_i^0H_j^\mp$, which appear in a nonminimal Higgs sector; here H_i^0 and H_j^0 are different neutral particles (for identical bosons the process $Z^0\to H^0H^0$ is forbidden by Bose statistics). The corresponding vertices are

$$A(Z^0\to H_i^0 + H_j^0) = i\frac{e}{\sin 2\theta_W} c_{ij}^{00}\varepsilon_\mu (p_i - p_j)_\mu,$$
$$A(W^+\to H_i^0 H_j^+) = \frac{e}{\sin 2\theta_W} c_{ij}^{0+}\varepsilon_\mu (p_i - p_j)_\mu.$$

Here ε_μ is the polarization of the Z^0 of W bosons, and c_{ij}^{00} and c_{ij}^{0+} are "mixing angles."

If charged Higgs bosons do exist in nature, then it would seem at first glance to be very tempting to seek them, like the standard H^0 boson, in associative production with gauge $W^\pm$ and Z^0 bosons. The standard $SU(2)\times U(1)$ theory, however, has no such vertex in the skeletal approximation, regardless of the number of Higgs doublets.[59–61] This vertex appears in a theory with a nonstandard Higgs sector which contains nondoublet as well as doublet Higgs fields. Since the existence of such fields violates the relation $\rho = M_W^2/M_Z^2\cos^2\theta_W = 1$, which holds quite well experimentally, their vacuum expectation values and therefore the ZWH vertex must be small[59–61]:

TABLE II. The new generation of installations with colliding e^+e^- beams (first phase).

Installation	Site	Startup	$\sqrt{s_{max}}$, GeV	L_{max}, 10^{31} cm$^{-2}\cdot$s^{-1}	σ (the beam energy spread at $\sqrt{s_{max}}$), GeV
SLC	Stanford, USA	Late 1986–early 1987	100	0.65	0.5
TRISTAN	KEK, Japan	1986 (late)	60	4	0.05
LEP	CERN, Geneva	1988 (late)	110	1.6	0.08

$$\xi \bar{g} m_W Z_\mu W_\mu^\pm H^\mp, \quad \xi \approx \sqrt{|1-\rho|} \leqslant 0{,}12. \qquad (3.11)$$

The ZWH vertex also arises in theories of the electroweak interaction based on the $SU(2)_L \times SU(2)_R \times U(1)$ gauge group, but here again the corresponding vertex is small[59–61]:

$$\xi \leqslant \frac{1}{4} \frac{m_{W_L}^2}{m_{W_R}^2} \leqslant 0.011 \quad \text{or} \quad \xi \leqslant 0.22 \frac{m_{Z_L}^2}{m_{Z_R}^2} \qquad (3.12)$$

(W_L and Z_L are ordinary vector bosons, while W_R and Z_R are heavy "right-handed" bosons).

Since experimentally we have $m_{H^+} > 16$ GeV (Sec. 4), the decay $Z^0 \to H^+ W^-$ is forbidden by mass, and the production of an $H^\pm W^\mp$ pair, in e^+e^- collisions, for example, could go through a virtual Z^0 boson. For the cross section for this process we have

$$\sigma(e^+e^- \to Z^0_{\mathrm{virt}} \to W^+H^-) = \xi^2 \frac{G_F^2 m_Z^4}{24\pi s} (1+v_e^2) \frac{1}{[1-(m_Z^2/s)]^2} \left(\frac{p_W^2}{s} + \frac{3m_W^2}{s}\right) \frac{2p_W}{\sqrt{s}} \qquad (3.13)$$

[see (3.9), (3.11), and (3.12) for the notation].

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 described here, and in general they have only certain specific symmetry limitations.

b) Weinberg's model of *CP* breaking[62]

In a model with several doublets of Higgs fields there may be a spontaneous breaking of *CP* symmetry. For a long time, this possibility appeared more attractive than the *CP* nonconservation in the Kobayashi-Maskawa model,[57] where the breaking was embodied in the initial Lagrangian, although indirect arguments based on cosmological considerations had also been advanced[64] against the idea of a spontaneous breaking of discrete symmetries. As was shown in Ref. 65, if there are at least three Higgs doublets complex vacuum expectation values of the Higgs fields may be energetically favorable with a *CP*-invariant Lagrangian. This circumstance does not, however, lead to a complex quark mixing matrix or thus a *CP* breaking *in* the interaction of vector bosons with quarks. On the other hand, the mixing matrix of the Higgs fields themselves does turn out to be complex, with the result that a *CP*-noninvariant interaction of Higgs bosons with quarks and leptons arises.

A characteristic property of the model is an extremely large dipole moment for the neutron, which may in fact contradict (Ref. 66, for example) the existing experimental limit[67]: $d_n < 3.6 \cdot 10^{-25}$. The theoretical prediction for ε'/ε also appears to be above the present experimental limit. The most specific feature of the model is the necessary existence of extremely light charged Higgs particles with masses up to 10–20 GeV; this prediction appears to be on the verge of contradicting experiment (Sec. 4).

c) Higgs particles and supersymmetry

An attractive possibility for a further development of the standard $SU(3)\times SU(2)\times U(1)$ theory is its " supersymmetry" generalization (Refs. 68 and 69, for example). In supersymmetry theories, particles with both half-integer spin and integer spin fall in the same multiplets. The numbers of fermion and boson degrees of freedom turn out to be balanced.

In supersymmetry theories it is possible to avoid one of the most serious difficulties of the ordinary theory of the electroweak interaction. In ordinary models (i.e., not supersymmetry models) the mass of the Higgs boson diverges quadratically. If we do not change the structure of the theory up to a Planck mass (10^{19} GeV), or even up to a grand unified mass ($\sim 10^{15}$ GeV), then we might expect huge masses for the Higgs bosons. The assumption that there is a "renormalization" term of the opposite sign in the initial Lagrangian puts the question in a slightly different form: How do two such large quantities cancel out exactly? The question becomes even more disturbing because, in a sense, the difference which remains after this cancellation of the "radiation" and "renormalization" masses of the Higgs boson is by no means zero; the Fermi constant of the weak interaction is expressed in terms of it. We must therefore have two quantities which cancel out to 25–32 significant digits, and the remainder determines all low-energy physics.

In supersymmetry theories the contributions of the fermion and boson degrees of freedom to the mass of the Higgs bosons cancel out in a natural way. Since this cancellation is of a group nature, it remains in force in all orders of perturbation theory. Since the generators of the supersymmetry do not carry internal quantum numbers, no pair of a fermion and a boson of ordinary particles falls in a single supermultiplet, and each of the ordinary particles must be complemented with a supersymmetry partner.

Interestingly, supersymmetry requires the existence of at least two Higgs doublets, of which one gives a mass to quarks with a charge of 2/3, while the other gives a mass to quarks with a charge of $-1/3$: It turns out to be impossible to write a supersymmetry Lagrangian in which the up and down quarks would acquire mass from the same Higgs field. A simple argument shows that it is impossible to get by with a single Higgs doublet in a supersymmetry theory: The superpartner of a Higgs doublet, which falls in the same chiral supermultiplet as this doublet, is a doublet of left-handed (or right-handed) fermions with the same internal quantum numbers ("shiggs"). The appearance of such fermions in the theory would obviously lead to a triangle anomaly, which could be cancelled out, however, if there were a second chiral supermultiplet with the opposite value of the weak hypercharge.

A detailed discussion of the structure of the Higgs sector and of the mass spectrum of Higgs particles in specific supermultiplet models goes beyond the scope of the present review. Here we will discuss only a general result which applies to a very broad class of supermultiplet models, strictly speaking, to nearly all models with a spontaneous or soft breaking of the supersymmetry.[70,71] It turns out that in these models the lightest neutral scalar Higgs boson must be lighter than the Z^0 boson. For the simplest case of two Higgs doublets the following spectrum of Higgs particles arises (in order of increasing mass): a light neutral scalar boson, a neutral pseudoscalar boson, a heavy neutral scalar boson, and a charged Higgs boson. The heavy neutral scalar boson is heavier than the Z^0 boson, and the square of the mass of the charged Higgs particle is equal to the sum of the squares of the masses of the pseudoscalar boson and the W boson.

In general, supersymmetry, by establishing an equivalence between the fermion and boson degrees of freedom, is clearly a weighty argument in favor of the existence of elementary Higgs fields.

d) Scalar particles and technicolor

Another interesting possibility for avoiding the problem of a quadratic divergence in the mass of scalar particles is to assume a dynamic spontaneous breaking of the gauge $SU(2)\times U(1)$ symmetry similar to the breaking of chiral symmetry in quantum chromodynamics. Such models have been labeled "technicolor" theories (see Ref. 72, for example, for a review of the work in this direction).

In the technicolor model the elementary Higgs bosons are replaced by composite particles constructed from "techniquarks" or "technileptons." These "technifermions" have a strong "technicolor" interaction with a small confinement radius which gives rise to various bound states. Three composite massless Goldstone degrees of freedom go into weighting the $W^{\pm}$ and Z^0 bosons; the masses of the other "technihadrons" are expected to be of the order of 1 TeV. The bound states of techniquarks and technileptons might, however, also include some lighter entities, the so-called pseudo-Goldstone technicolor bosons.[73] The existence of these particles is associated with the global symmetry of the technicolor interaction under unitary chiral transformations of the various technileptons and techniquarks. The spontaneous breaking of this symmetry gives rise to pseudoscalar Goldstone bosons which are not "eaten" by gauge bosons and which acquire a relatively small mass through other interactions. Among the lightest pseudo-Goldstone technicolor bosons expected in the simplest models are[73] two neutral bosons P^0 with a mass $\lesssim 3$ GeV and a pair of charged bosons $P^{\pm}$ with a mass between 5 and 15 GeV.

The properties of these pseudo-Goldstone technicolor bosons are similar to those of ordinary light Higgs particles primarily because their coupling constants with the various quarks and leptons are proportional to fermion masses. It is thus difficult to distinguish pseudo-Goldstone technicolor bosons from elementary Higgs particles. For example, this prediction of the existence of $P^{\pm}$ with a mass $\lesssim 15$ GeV verges on contradicting experiment (Sec. 4). One distinction between pseudo-Goldstone technicolor bosons and elementary Higgs bosons is that there are no ZZP^0 or $W^+W^-P^0$ vertices for the former at the skeletal level. In particular, they could thus not undergo an associative production with Z and W bosons.

e) The axion

The axion was introduced[74,75] to solve the problem of the natural conservation of *CP* parity in the strong interaction—the so-called θ problem.[76] To review the theory of the axion would be to go beyond the scope of the present review (see Ref. 77, for example). Here we will simply describe the experimental status of the axion.

The "standard" axion is a pseudoscalar Higgs particle with a mass of the order of a few hundred keV (strictly speaking, all that the theory tells us is that this mass is greater than 150 keV, but the absence in experiments[78] of $a\to e^+e^-$ decays forces us to assume $m_a < 2m_e$) and a lifetime[74,79,80]

$$\tau\,(a\to 2\gamma)\approx(0.8\text{ s})\left(\frac{100\text{ keV}}{m_a}\right)^5.$$

The interaction of the axion with nucleons is reminiscent of the interaction of the π^0, but with a coupling constant smaller by a factor $\sim 10^3$.

The existence of the standard axion may now be regarded as an essentially settled matter: The absence of 2γ events in the decay of axions near the working reactor at Jülich[81] completely contradicts the predictions of the theory (see also the earlier studies[82] and the work by the collaboration[83] between the Joint Institute for Nuclear Research and the Scientific-Research Institute of Nuclear Physics at Moscow State University). Further evidence against the standard axion is the experimental absence of the decays $J/\psi\to a\gamma$[$B(J/\psi\to a\gamma) < 1.4\cdot 10^{-5}$ (90% c.l.)[84]] and $\Upsilon\to a\gamma$, $\Upsilon'\to a\gamma$[$B(\Upsilon\to a\gamma) < 3\cdot 10^{-4}$ (90% c.l.; see Ref. 85 and also Ref. 86]. Here we have a product $P = B(J/\psi\to a\gamma)\,B(\Upsilon\to a\gamma) < 4.2\cdot 10^{-9}$, in substantial contradiction of the theoretical prediction $P=(1.6\pm 0.3)\cdot 10^{-8}$ (the prediction for the product P, in contrast with the predictions for each of the decays separately, does not depend on the unknown parameter x: the ratio of the vacuum expectation values of the Higgs fields). A sensitive test of the existence of the axion is the decay $K^+\to\pi^+ + a$. This decay has been analyzed in several theoretical papers.[74,80,87,88] The predictions regarding $B(K^+\to\pi^+a)$ have ranged from 10^{-5} to 10^{-8}. An extremely high limit, $B(K^+\to\pi a) > 3.5\cdot 10^{-5}$, was found in Ref. 89, but the estimates there are not entirely convincing. On the other hand, decays through a c-quark loop yield[90] $B(K^+\to\pi^+a) = 2\cdot 10^{-5}x^2$. The new experimental limit[91] $B(K^+\to\pi^+a) < 4.8\cdot 10^{-8}$ is hardly compatible with the existence of the standard axion.

It has been suggested in several studies that it might be possible to retain the axion in the form required for solving the θ problem while simultaneously suppressing its interaction with quarks and leptons (an "invisible axion"). This refinement was made in Ref. 92 at the cost of introducing an additional heavy quark, while in Refs. 93–96 an additional complex field Φ which is a singlet under the Weinberg-Salam group was introduced.

A distinctive feature of the invisible axion is that it consists nearly entirely of the singlet field Φ, containing only a small admixture λ of ordinary doublet fields φ: $\lambda\sim\langle\varphi\rangle/\langle\Phi\rangle$, $\langle\varphi\rangle\ll\langle\Phi\rangle$. This circumstance weakens the interaction of the axion with quarks and leptons by a factor of λ and reduces its mass by roughly the same factor. In the SU(5) grand unification model we have $\langle\varphi\rangle \sim 100$ GeV and $\langle\Phi\rangle \sim 10^{15}$ GeV, i.e., $\lambda \sim 10^{-13}$ (Ref. 95).

It was subsequently shown that too large a value, $\langle\Phi\rangle \gtrsim 10^{12}$–$10^{13}$ GeV, leads to an energy density of axions in the modern epoch which is unacceptable from the cosmological standpoint.[97] The most popular range of values of $\langle\Phi\rangle$ is $10^8 \lesssim \langle\Phi\rangle \lesssim 10^{13}$ GeV. The lower boundary here arises from the requirement that the emission of axions by red giants not affect their evolution.[98–100] Sikivie[101] has pointed out that such invisible axions might be observable experimentally from their conversion into γ rays in a strong nonuniform magnetic field.

f) The arion

The arion is a strictly massless Goldstone boson which is associated with spontaneous breaking of the exact chiral symmetry.[102] Here we will simply review the experiments which have placed limitations on its interactions.

The interaction of the arion with quarks and leptons is described by

$$L = x_f \frac{m_f}{v} (\bar{f} i\gamma_5 f)\,\alpha, \quad v = (G_F\sqrt{2})^{-1/2} = 246\ \text{GeV};$$

here m_f is the mass of the fermion, f and α are fermion and arion fields, and the dimensionless parameter x_f is associated with the ratio of the various vacuum expectation values. This ratio is, generally speaking, of the order of unity. A value $x_f \ll 1$ is, however, possible.

The exchange of an arion leads to long-range spin-spin forces between quarks and leptons similar to a very weak magnetic interaction of spins. An attempt might be made to detect this long-range force in macroscopic experiments.[102,103]

The experimental absence of the axion decays $J/\psi \to a + \gamma$ (Ref. 84), $\Upsilon \to a + \gamma$ (Ref. 85), and $K^+ \to \pi^+ + a$ (Ref. 91), which we have already mentioned, also means that there are no analogous arion decays, since the vanishing mass of the axion is unimportant in these experiments. It can therefore be asserted that the conditions $x_c < 0.6$ (Ref. 84) and $x_b < 1$ (Ref. 85) hold for the arion. As for the decay $K^+ \to \pi^+ + \alpha$, we note that an estimate[90] of the emission of an axion (or an arion) by a c quark, $B(K^+ \to \pi^+\alpha) < 2\cdot 10^{-5} x_c^2$, leads to the limitation $x_c < 0.05$. On the other hand, theory asserts[102] $x_u = x_c = \ldots = -x_d = -x_s = \ldots x_q$.

A limitation on the direct interaction of the arion with light u and d quarks follows from the experimental absence of an anomalous splitting of F levels in ortho-hydrogen molecules.[104] The limitation is not very severe: $x_u < 3.5$. From the agreement of the anomalous magnetic moments of the muon and the electron with the predictions of quantum electrodynamics[105] we easily find $x_\mu < 3.6$ and $x_e < 100$.

Much more severe restrictions arise from astrophysical considerations.[98–100] Since an arion interacts only slightly with matter, it can freely escape from stars, so that stars would lose energy very rapidly. The requirement that the arion luminosity of the sun not exceed the photon luminosity leads to the condition $x_e < 10^{-3}$. An even more stringent limitation emerges from an examination of the evolution of red giants: $x_e < 10^{-6}$–10^{-7}. This estimate, however, is more model-dependent. The astrophysical data are thus crucial to the possible existence of arions.

In a special experimental search[106] for a long-range arion interaction a study was made of the precession of the nuclear spins of the mercury isotopes Hg^{199} and Hg^{201} in an arion field whose source consisted of the oriented spins of a ferromagnetic shield. This experiment yielded a limitation on the product $x_e x_u$: $|x_e x_u| < 2.5\cdot 10^{-3}$ (unfortunately, this limitation incorporates a theoretical uncertainty related to the complexities of nuclear calculations for mercury isotopes). There are plans to improve this limitation by at least one or two orders of magnitude.

g) The familons

A theoretical mechanism which gives rise to massless Goldstone bosons (or superlight pseudo-Goldstone bosons, *of the axion type*), *during whose emission there are changes* in the fermion flavors (strangeness, muon charge, etc.), was pointed out in Refs. 107 and 108. Such bosons were labeled "familons" in Ref. 107. Familons arise if the theory contains a "horizontal" symmetry: a strongly broken gauge symmetry between generations of quarks and leptons. Another necessary condition is that the Lagrangian have the additional global U(1) symmetry of the general chiral transformation of the fermions of various generations. Since the breaking of the horizontal group should occur at very short range,[109] the familons, like the invisible axion, interact only slightly with matter and have a very small mass (in certain versions of the theory they are strictly massless).

An example of a theory with a familon within the framework of the realistic $SU(5)\times SU(3)_h$ model[110] was analyzed in Ref. 108. The interaction of a familon α with quarks and leptons is described in this model by the Lagrangian

$$L = \frac{\sqrt{2m_d m_s}}{\langle\eta\rangle}\,\alpha\,(\bar{d}s + \bar{s}d) + \frac{\sqrt{2m_e m_\mu}}{\langle\eta\rangle}\,\alpha\,(\bar{\mu}e + \bar{e}\mu);$$

here $\langle\eta\rangle$ is the vacuum expectation value of one of the Higgs fields responsible for the breaking of the horizontal group. This Lagrangian leads to the decays $\mu \to e + \alpha$ and $K^+ \to \pi^+ + \alpha$, for which the probabilities are

$$\Gamma(\mu \to e\alpha) = \frac{1}{8\pi}\,\frac{m_e m_\mu^2}{\langle\eta\rangle^2}, \qquad \Gamma(K^+ \to \pi^+ + \alpha) \approx \frac{1}{8\pi}\,\frac{m_d m_K^3}{\langle\eta\rangle^2 m_s}.$$

We thus have an experimental limit on the vacuum expectation value $\langle\eta\rangle$. A more stringent limit comes from the experimental absence of this decay $K^+ \to \pi^+\alpha$: $\langle\eta\rangle > 10^{10}$ GeV. This value is far larger than the typical masses which emerge from the standard limitations on flavor-changing neutral currents.[109] If familons did in fact exist, the decays $K^+ \to \pi^+\alpha$ and $\mu \to e\alpha$ would be incomparably easier to observe than any ordinary effects of flavor-changing neutral currents ($\mu \to 3e$, $\mu \to e\gamma$, $K_L \to e^\pm \mu^\mp$, $K^+ \to \pi^+ e^\pm \mu^\mp$ etc.).

h) The majoron

The possibility that neutrinos are not strictly massless particles is attracting increasing interest. Undoubtedly the strongest piece of evidence in favor of this possibility is an experiment carried out at the Institute of Theoretical and Experimental Physics to measure the edge of the tritium β

spectrum.[111] There is an active discussion in the literature of models in which the appearance of a Majorana mass of neutrinos is linked with spontaneous breaking of lepton-number conservation and thus the appearance of a corresponding massless Goldstone boson: the "majoron."[112,113]

A neutrino which is a right-handed singlet under the Weinberg-Salam group is introduced in the model of Ref. 112. By virtue of an ordinary Yukawa coupling with a doublet Higgs field, the neutrino acquires a "Dirac" mass m, which is generally of the order of ordinary lepton masses. It is furthermore assumed that the right-handed neutrino has a large Majorana mass M by virtue of a Yukawa coupling with a singlet Higgs field. In this case the physical particles with definite masses are two Majorana neutrinos with masses M and m^2/M. The heavy neutrino consists almost entirely of a singlet right-handed neutrino and a singlet left-handed antineutrino, while the light Majorana neutrino consists essentially of an ordinary left-handed doublet neutrino and a right-handed doublet antineutrino.

The formation of the vacuum expectation value of the singlet scalar field gives rise to a spontaneous breaking of lepton number and to the appearance of a massless Goldstone boson: the majoron. The majoron is coupled comparatively strongly with a heavy neutrino (the coupling constant can be of the order of the ordinary Yukawa coupling constants), and it is coupled considerably less strongly [by a factor of $(m/M)^2$] with an ordinary neutrino.

In addition there is a nondiagonal coupling of intermediate strength of a majoron with both neutrinos, which gives rise to a rapid decay of the heavy neutrino into a light neutrino and a majoron. The interaction of the majoron with quarks and leptons arises only in the single-loop approximation and is essentially unobservable. One manifestation of the existence of the majoron in this model might be the possibility of the decay of the ν_μ (and the ν_τ) into a ν_e plus a majoron. If this decay occurred in a time shorter than the age of the universe, the mass of ν_μ (or ν_τ) would not be restricted by the known astrophysical limit (the sum of the masses of all stable neutrinos is < 30eV). Actually, however, for reasonable values of the parameters the lifetime of the ν_μ turns out to be longer than the age of the universe.[114]

Another majoron model, in which the neutrino acquires mass without the introduction of additional fermion degrees of freedom, was offered by Gelmini and Roncadelli.[113] In this model a triplet of Higgs fields with a small vacuum expectation value of the neutral component gives rise to a small Majorana mass for the neutrino. The majoron consists almost entirely of a neutral component of a triplet field which is coupled only with neutrinos, and it contains only a small admixture of a doublet field. It thus interacts only slightly with quarks and leptons.

Possible experimental manifestations of the existence of such a majoron have been discussed in several papers.[115–119] It follows from an estimate of the energy radiated by red giants that the triplet vacuum expectation value is less than 100 keV, and a comparison of the predictions of the model with experiments on double β decay reveals that the constant of the interaction of the majoron with an electron is $< 10^{-3}$, and the mass is limited by[115] $m_{\nu_e} < 15$ eV. There have also been studies[116–119] of decays of K and π mesons involving a majoron,[116] corrections to μ decay,[117] the decay[118] $\mu \to e\gamma$, and possible hydrogen $\rightleftarrows$ antihydrogen oscillations.[119]

On the whole, we now have a rather large literature on both astrophysical and possible laboratory manifestations of the majoron. There is, however, the real difficulty that it is not possible to specify any definite phenomenon (or phenomena) which might yield an unambiguous answer to the question of its existence.

4. EXPERIMENTAL LIMITATIONS ON THE EXISTENCE OF HIGGS BOSONS

The present experimental lower limits on the mass of the standard Higgs boson, H^0 are considerably weaker than theoretical conditions (1.13) and (1.13a). Let us examine these limitations.

For a scalar boson which interacts with fermions in accordance with (1.4), an analysis of neutron scattering by atomic electrons yields the limitation[120] $m_H > 0.6$ MeV. A Higgs boson with a mass less than 13 MeV is incompatible with the measured angular distributions in low-energy neutron-nucleus scattering.[121] The negative results of experimental searches for $0^+ \to 0^+$ nuclear transitions of ^{16}O (6.05 MeV) and ^4He (20.2 MeV) to the corresponding ground states, accompanied by the emission of H^0 and the subsequent decay $H^0 \to e^+e^-$, is evidence that the Higgs boson cannot have a mass in the interval 1.030 MeV $< m_H < 5.84$ MeV [^{16}O (6.05 MeV); Refs. 122 and 120][5)] and 2.8 MeV $\leqslant m_H < 11.5$ MeV [^4He (20.2 MeV).[124] The absence of anomalies from the x-ray spectra of μ-atoms leads to roughly the same lower limit[125] on m_H.

A search for $J/\psi \to \gamma a$ decays (a is an axion) carried out by the Crystal Ball group on SPEAR (Stanford Linear Accelerator Center) yields the limitation $B(J/\psi \to \gamma a) < 1.4 \times 10^{-5}$ (90% c.l.).[84] This limitation applies not only to the axion but also to any other sufficiently light particle which does not manage to decay inside the apparatus. For a standard Higgs boson, this condition means[84] $m_H < 50$ MeV. The experimental limit found for $B(J/\psi \to \gamma H^0)$ is lower than the theoretical value (2.18) by a factor of at least three. It thus follows that[84] $m_H > 50$ MeV.

The absence of $K^+ \to \pi^+ + H^0$ decays imposes the limit $m_H > 350$ MeV. The width of this decay is relatively large[6]:

$$\frac{\Gamma(K^+ \to \pi^+ + H^0)}{\Gamma(K^+ \to \text{all})} = 2.1 \cdot 10^{-4} \cdot \sqrt{1 - 0.18\,\frac{m_H^2 - m_\pi^2}{m_\pi^2}}\,\sqrt{1 - 0.05\,\frac{m_H^2 - m_\pi^2}{m_\pi^2}}. \tag{4.1}$$

[5)]It should be noted, however, that the analysis of Refs. 120–122 was based on the assumption that the constant of the interaction of H^0 with nucleons, g_{H^0NN}, is equal to m_N/v. Actually, the direct interacion of H^0 with light quarks is slight, and the H^0NN vertex is determined by interaction (1.5a) of the H^0 with gluons. Here we have [123] (see also Ref. 5)

$$g_{H^0NN} = \frac{2n_h}{27}\,\frac{m_N}{v} = \frac{70\ \text{MeV}}{v}\,n_h$$

(n_h is the number of heavy quarks). This circumstance eases the limitations of Refs. 120–122 slightly.[14,124]

If the H^0 does not decay inside the apparatus, the process $K^+ \to \pi^+ + H^0$ would simulate the decay $K^+ \to \pi^+ \nu\tilde{\nu}$, which is absent, at least at the level of $4.8 \cdot 10^{-8}$ (Ref. 91). Experimentally, the decay $K^+ \to \pi^+ e^+ e^-$ has a probability of about $2.7 \cdot 10^{-7}$, while that of $K^+ \to \pi^+ \mu^+ \mu^-$ is no more than $2.4 \cdot 10^{-6}$. On the other hand, a Higgs boson would decay primarily into an e^+e^- or $\mu^+\mu^-$ pair (in the interval $2m_\pi < m_H < 350$ MeV the probability for the decay of H^0 by the $\mu^+\mu^-$ mode is higher than that for the $\pi\pi$ mode).

An H^0 boson with a slightly larger mass (up to 408 MeV) might be observed in the decay $\eta' \to \eta H^0 \to \eta \mu^+ \mu^-$, but the existing experimental limitation[126] $B(\eta' \to \eta \mu^+ \mu^-) < 1.5 \times 10^{-5}$ (90% c.l.) apparently cannot be taken as reliable evidence for the absence of a standard Higgs boson with a mass greater than 350 MeV, if we take into account the possible uncertainty in the theoretical calculations of $\Gamma(\eta' \to \eta H^0)$ and the existing experimental uncertainty regarding the total width of the η' meson.

In order to use the radiative decays of J/ψ to search for a heavier Higgs boson, the Crystal Ball group mentioned earlier also studied the decays $J/\psi \to \mu^+\mu^-\gamma$ (Ref. 84) and $J/\psi \to \pi^0\pi^0\gamma$ (Ref. 127) and the distribution of events in the invariant masses of the $\mu^+\mu^-$ and $\pi^0\pi^0$ pairs. Analysis of the data revealed that at $B(H^0 \to \mu^+\mu^-) > 0.75$ the interval 400 MeV $< m_H < 1$ GeV is forbidden,[84] while at $B(H^0 \to \pi\pi) > 0.7$ the interval 500 MeV $< m_H < 1$ GeV is forbidden.[127]

Consequently, the present experimental limitation is $m_H > 350$ MeV.

Fans of light Higgs particles received some good news in the summer of 1983 when the MARK-III group announced the discovery of a new narrow neutral resonance,[128] $\xi(2.2)$ with $m_\xi = 2.22 \pm 0.01$ GeV, $\Gamma_\xi < 40$ MeV, in agreement with the experimental resolution. In radiative decays of J/ψ, the decay channels $\xi \to K^+K^-$ and $\xi \to K_S^0 K_S^0$ were observed, with

$$B(J/\psi \to \gamma\xi)\, B(\xi \to K^+K^-) = (5{,}8 \pm 1{,}8 \pm 1{,}5) \cdot 10^{-5}. \tag{4.2}$$

The unusually small width of the ξ particle and the observed $K\bar{K}$ decay mode were discussed as an argument that the ξ particle had to be tested before it could be identified as belonging to the Higgs sector; see Ref. 129, for example. We should emphasize at the outset that even if experiment were to establish that ξ is a scalar particle (the decay $\xi \to K_S K_S$ means only that $J_\xi^{PC} = 0^{++}, 2^{++}, \ldots$) this particle cannot be regarded as a standard-model Higgs boson with a single doublet. Aside from the fact that its mass is well below the limit in (1.13), the probability for the transition $J/\psi \to \xi\gamma \to K^+K^-\gamma$ in (4.2) is considerably higher than that expected for the standard H^0 boson. It follows from (2.18) that with $m_H \approx 2.2$ GeV we would have $B(J/\psi \to \gamma H^0) \approx (3.1 \pm 0.5) \cdot 10^{-5}$. For comparison with (4.2), this value should also be multiplied by $B(H^0 \to K^+K^-)$. Under the assumption that the $\bar{s}s$ mode is dominant, we would expect that at least the following quasi-two-particle final states would be realized: K^+K^-, $K^0\bar{K}^0$, $K^*\bar{K}^*$, $K\bar{Q} + \bar{K}Q$. Here we have the conservative estimate $B(H^0 \to K^+K^-) \lesssim 0.1$–$0.2$. Consequently, the theoretical prediction for the standard H^0 with a mass of 2.2 GeV is about an order of magnitude, lower than the observed value in (4.2).

The arguments above are literally inapplicable in the case of a more complicated Higgs sector, containing several doublets of scalar fields (Sec. 3), in which case light scalar particles would also be possible. The expected value of $B(J/\psi \to \gamma H^0)$ might increase in this case because of a decrease in the vacuum expectation values and thus an increase in the corresponding Yukawa constants [Sec. 3; expression (3.4)]. Nevertheless, even here the hypothesis that ξ is a Higgs boson is not stimulating any special optimism.

In particular, there are caution flags such as the experimental absence of signals corresponding to $J/\psi \to \gamma\xi \to \gamma\mu^+\mu^-$ and Υ, $\Upsilon' \to \gamma\xi$: $B(\Upsilon \to \gamma\xi)\, B(\xi \to K^+K^-) < 2 \cdot 10^{-4}$ (90% c.l.) and $B(\Upsilon' \to \gamma\xi)\, B(\xi \to K^+K^-) < 9 \cdot 10^{-5}$ (90% c.l.).[130] For a Higgs boson we might expect that the decay $\xi \to \mu^+\mu^-$ would be completely observable $[B(\xi \to \mu^+\mu^-) \sim 0.1]$ and that the transitions $\Upsilon, \Upsilon' \to \gamma H^0$ would be further enhanced by a factor $\sim (m_b/m_c)^2$. In any case, the fact that these decays are not observed places a serious restriction on the class of models with a nonminimal Higgs sector, in which ξ could be a Higgs particle.

If future measurements confirm the hypothesis $J_\xi^{PC} = 0^{++}$ and provide clear evidence against the interpretation of ξ as a hadron resonance (this possibility seems unlikely), a necessary test of the Higgs nature of this state would be a search for its production along with $W^\pm$ or Z^0 bosons. Since the ξ particle has a small mass ($m\xi < \Gamma_W, \Gamma_Z$), the corresponding cross sections for its production might be large (Sec. 2). In particular, for the $Sp\bar{p}S$ collider we could expect $\sigma_{\xi W} \approx 5 \cdot 10^{-2}$ mb.

Attempts to observe the Higgs boson have already been undertaken at the Υ resonance, through measurements of the spectrum of inclusive γ rays from its decay $\Upsilon \to \gamma + \text{all}$. An experiment by the CUSB group[131] on the CESR (Cornell) yielded the limitation $B(\Upsilon + \gamma + H^0) < (1$–$2.5) \cdot 10^{-3}$(90% c.l.). This limitation is about an order of magnitude weaker than that required for a standard H^0 boson with $m_H \lesssim 8$ GeV [see (2.18)].

Let us briefly review the existing experimental limitations on the existence of charged Higgs bosons.

We know that the decays of τ leptons and charmed hadrons are described well by the standard theory. This circumstance imposes a limit $m_{H^+} > m_\tau$. Limitations on the masses of charged Higgs bosons also follow from an analysis of data on the decays of charmed B mesons ($B^- = b\bar{u}$, $B^0 = b\bar{d}$)[132] and on the production in e^+e^- annihilation of noncollinear hadron jets and/or direct hard leptons.[133] The decays of B mesons rule out the existence of H^+ with masses $m_\tau < m_{H^+} < m_b - m_c$. The various characteristics of the B decays agree well with the standard model for the transition $b \to c + W^-$, and they completely refute a dominant role of the cascades

$$\begin{array}{l} b \to qH^- \\ \quad \lfloor\!\longrightarrow \tau^-\nu_\tau,\ \bar{c}s. \end{array} \tag{4.3}$$

Corresponding experiments have been carried out by the CLEO group on the CESR.[132] The most important experiments were measurements of the yields of inclusive e and μ and of the fraction of the energy carried off by charged hadrons in the decay of the B meson. Cascades (4.3) would have substantially reduced these characteristics in comparison with the standard model for the decay of the b quark. The limitations which have been found also apply to charged technicolor bosons.

The observation of the t quark with the standard decay model $t \to b + W^+$ rules out the existence of a Higgs boson with a mass $m_{H^+} < m_t - m_b$. The confirmation of the value $m_t = 40$ GeV found in the experiment on the Sp$\bar{p}$S collider and of the standard properties of the t decay would impose a new restriction $m_{H^+} > 35$ GeV.

In e^+e^- annihilation at the energies presently atainable, $\sqrt{s} \lesssim 46\ \mathrm{GeV} < 2m_t$, the binary production and subsequent decays of charged Higgs bosons might give rise to the reactions

$$e^+e^- \to H^+H^- \left| \begin{array}{l} \to \tau^+\tau^-\nu_\tau\bar{\nu}_\tau, \\ \to c\bar{s}\tau^-\bar{\nu}_\tau,\ \bar{c}s\tau^+\nu_\tau, \\ \to c\bar{s}c\bar{s}, \end{array} \right. \tag{4.4}$$

which would have the following experimental manifestations: 1) noncollinear pairs of τ leptons, 2) events with a hard lepton jet and a hadron jet, 3) four-jet hadron events corresponding to the ($\bar{c}$s) ($\bar{c}$s) mode.

A search for such events has been carried out by several experimental groups on PETRA (Hamburg, West Germany) and PEP (Stanford).[134] A detailed analysis of the data rules out the possible existence of $H^\pm$ with masses in the interval $m_{H^+} \sim 5$–16 GeV.

THE OUTLOOK

High-energy physics has been complicated in recent years by the circumstance that the startup of essentially each new accelerator has been accompanied by the discovery of some fundamental new entity or phenomenon. An important landmark along this route was the 1983 discovery of the $W^\pm$ and Z^0 bosons on the Sp$\bar{p}$S collider at CERN,[1,2] which has shared in this good fortune.

The best candidates for the next sensations in particle physics are the new experiments on the Sp$\bar{p}$S collider and the startup of some tough competitors: the accelerators with colliding e^+e^- beams of the next generation (see, for example, Ref. 135 and Table II). One of the foremost problems for all these installations, of course, is a critical experimental test of the possible existence of Higgs bosons. It is being suggested that research on the Higgs sector be continued on the ultrahigh-energy accelerators of the next generation, with colliding e^+e^-, $p\bar{p}$, and pp beams, which are presently being developed. Here is a list of the most important of these projects.

The possibility of further increases in the energy and luminosity L in future phases of LEP is being discussed. The outlook for substantially increasing the energy of colliding e^+e^- beams depend on the successful development of the method of linear colliding beams. At Novosibirsk, for example, the VLÉPP project with $\sqrt{s} = 300$–1000 GeV is being developed.

Plans call for the 1986 startup of the Tevatron—a $p\bar{p}$ collider with $\sqrt{s} = 2$ TeV and $L \approx 10^{30}$–10^{31} cm^{-2}s^{-1}—at Fermilab. The LHC, a $p\bar{p}$ collider with $\sqrt{s} = 10$ TeV and $L = 10^{31}$ cm^{-2}s^{-1}, is expected to be constructed in the LEP Tunnel in the early 1990s. Plans call for the UNK accelerator complex (Serpukhov), already constructed, to switch to operation with colliding pp beams with $\sqrt{s} = 6$ TeV and $L = 10^{32}$ cm^{-2}s^{-1} in the future. Finally, the SSC collider with $\sqrt{s} = 40$ TeV and a luminosity $L \approx 10^{33}$ cm^{-2}s^{-1} (pp) or $L = 10^{32}$ cm^{-2}s^{-1} ($p\bar{p}$) has been approved in principle in the USA. This collider can be expected to come on line in the early 1990s.

Because of the particular features of processes involving the production and subsequent decay of Higgs bosons, the experimental groups which intend to join the hunt for these bosons are developing modifications of existing detectors and are inventing new instruments in order to create the most favorable conditions for observing these particles. The experiments which have been planned include calorimetric measurements and a good identification of hard leptons for detecting the decays of $W^\pm$ and Z^0 bosons and heavy quarks. Plans call for the use of vertex detectors to observe the ranges of those heavy particles with lifetimes of 10^{-13}–10^{-12} s which are inseparable companions of Higgs bosons (Sec. 1). The UA1 and UA2 groups, for example, which achieved the long-awaited discovery of $W^\pm$ and Z^0, are modernizing their detectors to reach new frontiers in particle physics. Let us take a brief look at the prospects from the standpoint of the search for Higgs bosons.

The total energy of the Sp$\bar{p}$S collider is expected to increase from the present $\sqrt{s} = 2 \times 270$ GeV to $\sqrt{s} = 2 \times 310$ (330) GeV. The average luminosity of this collider is to be increased to $L \sim (1$–$2)\ 10^{30}$ cm^{-2}s^{-1}. It is assumed that over 3 yrs of operation of the UA1 detector an integrated luminosity $L_{\mathrm{int}} \sim 10^4$ nb^{-1} will be achieved.[44]

Plans call for a search for a Higgs boson H^0 with a mass up to 40 GeV on the basis of the associative production of this boson with $W^\pm$ (Sec. 2). The lepton modes and quark modes of the decay of $W^\pm$ are to be identified, and H^0 is also to be detected on the basis of the heavy products of its decay. Monte Carlo calculations[44] of the expected number of events corresponding to $W^\pm H^0$ production at $L_{\mathrm{int}} = 10^4$ nb^{-1}, $\sqrt{s} = 660$ GeV, and various values of m_H are listed in Table III.

TABLE III. Expected number of events corresponding to the process $\bar{p}p \to H^0 + W^\pm$ + all at $\sqrt{s} = 660$ GeV with $L_{\mathrm{int}} = 10^4$ nb^{-1}. According to calculations based on the actual capabilities of the UA1 installation[44].

m_H, GeV	$H^0W^\pm$ production with detection of the decay $W \to e_\nu^{(-)}$ or $W \to \mu_\nu^{(-)}$	$H^0W^\pm$ production with detection of the decay $W \to 2$ jets
10.4	27	108
20	10.8	43
30	3.6	14.5
40	1.6	6.5

The modification of the detector and the accelerator will take 2–3 yrs, and another 3 yrs or so will be required in order to acquire a sufficient statistical base. The Sp$\bar{p}$S thus has a good chance to carry out the first search for H^0 with $m_H < 40$ GeV, before the corresponding SLC and LEP programs are implemented (Table II). A study of the region of larger masses m_H on the Sp$\bar{p}$S collider will require a substantial increase in its luminosity. If this increase proves impossible, the mass region $m_H > 40$ GeV will be studied first only in e^+e^- collisions or on pp ($\bar{p}$) colliders of the next generation.

What are the prospects for installations with colliding e^+e^- beams? The standard procedure described above (Secs. 3 and 4) for searching for charged Higgs bosons $H^{\pm}$ will be pursued at all available energies of e^+e^- annihilation, and we will not discuss this approach any further here. The choice of the process most favorable for a search for the neutral boson H^0 depends on whether toponium (T) is discovered in the near future and on precisely what its mass m_T is. Toponium with a mass of about 60 GeV, for example, would be a genuine gift to the TRISTAN installation. As was discussed in Sec. 2, toponium is an extremely convenient entity for searching for H^0 with $m_H < m_T$, on the basis of the decay $T \to \gamma = H^0$. The evidence for the production of H^0 in this decay would be monochromatic γ rays and a pair of heavy particles from the H^0 decay.

To illustrate the situation, let us estimate the expected number of events, assuming an average luminosity $L_{av} \approx 1 \cdot 10^{31}\ cm^{-2}s^{-1}$ and an energy spread $\sigma \approx 70$ MeV in the beams. Assuming a value $\Gamma(T \to \gamma_{virt} \to e^+e^-) \approx 5$ keV for the electromagnetic contribution to the lepton decay width of T, and calculating the cross section for e^+e^- annihilation at the peak of the T resonance by the standard procedure, allowing for radiation effects (Refs. 41, 42, and 52, for example), we find that at $m_T \approx 60$–80 GeV we would expect ~ 100–50 T production events per day. If the energy ($\sqrt{s}$) range below $2m_t$ is scanned at a step ~ 100 MeV, in accordance with conservative estimates,[136] 10^2 days would be required to detect a peak. At the value $B(T \to H^0 + \gamma) \approx 2 \cdot 10^{-2}$, the yield of Higgs bosons H^0 at the T peak should be observed at the level of one event per day.

If the searches for Higgs bosons before the beginning of experiments on the SLC and the LEP will turn out to be inconclusive, we can expect substantial progress in clarifying the question of the existence of these bosons through a detailed study o the reaction $e^+e^- \to H^0 + Z^0$ described above (Sec. 2). Analysis of the expected characteristics of the H^0 signal and of the background conditions (Refs. 134–136, for example) raises the hope that the Higgs boson will be discovered on the SLC or the LEP in the decay of Z^0 if its mass does not exceed ~ 50 GeV. Most of the hope for expanding the search H^0 to $m_H \sim 50$–100 GeV should apparently be pinned on the future phases of these installations, in which H^0Z^0 production can be studied above the threshold.

Table IV compares the yields of H^0 bosons corresponding to the various mechanisms for their production in e^+e^- annihilation at an integrated luminosity $L_{int} = 10^2\ nb^{-1}$ (see also Ref. 136). We should emphasize that the discovery of H^0 bosons with masses substantially greater than m_Z in e^+e^- collisions, like a study of other important questions in particle physics, will require accelerators with an adequately high luminosity. For the associative production of H^0Z^0, for example, the observation of a few events per day will require a luminosity $L = (2m_H/100\ \mathrm{GeV})^2 \times 10^{32}\ cm^{-2} \cdot s^{-1}$.

We will complete this section of the review with a brief comparison of the possibilities of searches for Higgs particles in e^+e^-, p$\bar{p}$, and pp collisions. The e^+e^- colliders have indisputable advantages from the standpoint of complete utilization of the entire collision energy, the best background conditions, and a small beam energy spread, which makes it possible to use "factories" of the vector bosons T and Z^0 (if $m_H < m_T \lesssim m_Z$). Another important advantage of these colliders is the fact that, in contrast with hadron colliders, the cross sections for the production of Higgs particles are generally a significant fraction of the total cross section for the production of hadrons. On the other hand, since the maximum energies of even the future phases of installations presently under construction will apparently not exceed[135] 200–250 GeV, the region of masses m_H above 150 GeV may prove to be inaccessible in practice for these installations.

Proton colliders might be able to search for Higgs particles over a broad mass interval (up to $m_H \lesssim 600$–700 GeV) through the gluon mechanism, in particular. Because of their substantially higher energies and higher luminosities (which are particularly important in the case of pp collisions), here it would be possible to achieve a rather high count rate of events corresponding to the production of Higgs bosons. However, an expansion of the m_H range accessible to measurements should result from a rapid increase in the energy $\sqrt{s}$. We emphasize that studies in the region $2m_t \lesssim m_H \lesssim 2m_W$ appear rather complicated because of the unfavorable background situation.

The e^+e^- installations are thus the best places to search for comparatively light H^0 bosons and to study their properties. In view of the dates at which they are expected to come on line, however (Table II), it may be that the Sp$\bar{p}$S collider at CERN will beat them and will furnish the first information on the existence of Higgs particles with $m_H < 40$ GeV.

The situation might change substantially if toponium with a mass $m_T < m_Z$ were to be discovered in the near fu-

TABLE IV. Yields of H^0 bosons in various reactions (events) in colliding e^+e^- beams at $L_{int} = 10^2\ nb^{-1*}$.

m_H, GeV	$e^+e^- \to T \to H^0 + \gamma$, $m_T = 75$ GeV, $\sigma = 70$ MeV	$e^+e^- \to Z^0 \to e^+e^-H^0$, $\sqrt{s} = m_Z$	$e^+e^- \to H^0 + Z^0 \to e^+e^-$, optimum energy
10	100	400	87
20	92	90	30
40	71	12	10
60	36	1	4.5
100	—	—	1
180	—	—	0.25

* $10^{31}\ cm^{-2} \cdot cm^{-1} = 0.86\ pb^{-1} \cdot day^{-1}$

ture, and a search for radiative decays $\Upsilon \to H^0 + \gamma$ were to be undertaken.

The outlook for the search for heavy Higgs particles in e^+e^- collisions depends on the highest energies which will be attained on the SLC and the LEP and also on developments in other projects, in particular, the VLÉPP. In any event, this search will be undertaken on the hadron colliders of the following generations.

Here we have attempted to map out a natural path for the search for Higgs bosons, on the basis of the standard assumptions about their properties. The optimist can of course always count on miracles. It is far from impossible that some new entity will unexpectedly pass the exam for honored membership in the Higgs sector and therefore radically revise the existing plans.

We wish to thank V. I. Zakharov, B. L. Ioffe, D. Klein, L. B. Okun', and M. A. Shifman for useful discussions.

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Translated by Dave Parsons