Signatures for charged-Higgs-boson production in e^+e^- collisions

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We investigate various signals for a charged-Higgs-boson search in e^+e^- annihilation. To facilitate our discussion, we introduce a simple parametrization of the charged-Higgs-boson couplings to fermions for three or more Higgs doublets in the Weinberg-Salam model.

I. INTRODUCTION

The minimal version of the Weinberg-Salam $SU(2)\times U(1)$ gauge model[1] of weak and electromagnetic interactions has proved to be remarkably successful in describing weak-interaction phenomenology. In this model there is just one Higgs scalar doublet before spontaneous symmetry breaking. Symmetry breaking leaves behind one neutral scalar Higgs particle, which remains to be observed. *CP* violation is introduced via the Kobayashi-Maskawa mechanism[2] with three or more doublets of quarks.

However, at very high temperatures, *CP* invariance is restored in this model when the quark masses vanish. If we believe the scenario for the production of a net baryon number density in the early universe, then according to grand unified models there must exist another origin of *CP* violation, which exists at temperatures $kT\sim 10^{14}$ GeV.[3] This type of *CP* violation, which we shall refer to as hard *CP* violation, can be incorporated into a grand unified model [e.g., the SU(5) model[4]] with the introduction of more Higgs multiplets. In the Weinberg-Salam model this means three or more Higgs doublets.[5] In this picture, *CP* is not an *a priori* symmetry of the Lagrangian; it arises as a consequence of the proliferation of Higgs scalar particles in the model, as originally observed by Weinberg.[5] This implies the existence of physical charged Higgs particles; they can have rather small masses (~ a few GeV) so that present and forthcoming experiments may be able to detect them.

In this note, we shall discuss the properties of charged Higgs particles and suggest some clear signals for their formation in e^+e^- annihilation experiments. The phenomenology of charged Higgs particles has been discussed extensively in the literature for the two-Higgs-doublet case.[6] However, with the proliferation of particles suggested here, the existing formalism for the description of Higgs-boson–fermion couplings[5] becomes rather cumbersome. We find it convenient to introduce a coupling matrix for the Higgs sector analogous to the Kobayashi-Maskawa matrix for the quark sector. This facilitates our discussion of charged-Higgs-boson phenomenology.

Recent measurements of R and jetlike signatures (e.g., thrust, sphericity) of the hadronic events at PETRA[7] from $E_{c.m.}=13$ to $\simeq 32$ GeV imply the absence of new charge-$\frac{2}{3}$ quarks and new heavy charged leptons in this energy range. This, in turn, forgoes the possibility of detecting charged Higgs bosons in the energy range covered by PETRA through the decays of heavy quarks or leptons with masses above the b-quark mass. Owing to the statistical errors in the measurement of R, however, the presence of charged Higgs particles with $m_H \leq 16$ GeV is not ruled out, since $\sigma(e^+e^-\to H^+H^-)$ contributes only $\frac{1}{4}$ of a unit to R. In view of the importance of searching for such objects, we have investigated ways to identify charged Higgs bosons assuming they can be pair-produced in e^+e^- annihilation directly via one-photon exchange.

The rest of the paper is organized as follows. In Sec. II, we discuss the formalism of the charged-Higgs-boson couplings to the fermions for three Higgs doublets. The generalization to more Higgs doublets is briefly mentioned. In Sec. III, we suggest methods to search for charged Higgs particles in e^+e^- annihilation. These complement the methods already suggested in the literature.[6] The results of this section are summarized in Table I.

II. WEINBERG-SALAM GAUGE MODEL WITH THREE HIGGS DOUBLETS

Let us consider the Weinberg-Salam model with three Higgs doublets

$$\phi_i = \begin{pmatrix} \phi_i^+ \\ \phi_i^0 \end{pmatrix}, \quad i=1,2,3. \tag{1}$$

To incorporate CP-violating effects into the Higgs sector, some of the couplings among the Higgs fields must be complex, as originally suggested by Weinberg.[5] After spontaneous symmetry breaking, there remain nine real scalar fields in the model, five of them neutral, two of them positively charged, and two of them negatively charged. Let us concentrate on the charged Higgs sector and refer to the mass eigenstates as $H_1^\pm$ and $H_2^\pm$. The original positively charged scalar fields, ϕ_1^+, ϕ_2^+, and ϕ_3^+, can be related to the Higgs Fields H_1^+ and H_2^+ and the charged Goldstone boson ϕ_W^+, which is absorbed into the W^+ by the unitary transformation

$$\begin{pmatrix} \phi_1^+ \\ \phi_2^+ \\ \phi_3^+ \end{pmatrix} = Y \begin{pmatrix} \phi_W^+ \\ H_1^+ \\ H_2^+ \end{pmatrix}. \tag{2}$$

The unitary Y matrix is analogous to the Kobayashi-Maskawa matrix. In general for an $n\times n$ unitary matrix, there are $(n^2-n)/2$ real parameters and $(n^2+n)/2$ phases. To ensure that flavor conservation is not violated by the ϕ_i^0 exchanges, we require that the Lagrangian be invariant under reflections of the ϕ_i (and the appropriate right-handed fermions) taken one at a time. (We shall come back to this point later.) Then, among the $(n^2+n)/2$ phases, n of them can be rotated away by redefining the ϕ_j by $\phi_j \to e^{i\Psi_j}\phi_j$ and $n-1$ of them can be rotated away by redefining the H_j by $H_j \to e^{i\chi_j}H_j$. Hence, there are $(n^2-3n+2)/2$ phases left behind before coupling the Higgs to the fermions. Since the fermion phases are all fixed, there is no further removal of phases in the Y matrix. This means that the Y matrix has $(n^2-n)/2$ real parameters, which can be expressed in terms of angles, and $(n^2-3n+2)/2$ phases, which give CP-violating effects to the fermion-Higgs sector.

For the three-doublet case, we have

$$Y = \begin{pmatrix} c_1' & s_1'c_3' & s_1's_3' \\ -s_1'c_2' & c_1'c_2'c_3'+s_2's_3'e^{i\delta'} & c_1'c_2's_3'-s_2'c_3'e^{i\delta'} \\ -s_1's_2' & c_1's_2'c_3'-c_2's_3'e^{i\delta'} & c_1's_2's_3'+c_2'c_3'e^{i\delta'} \end{pmatrix}_Y, \tag{3}$$

where c_i' (s_i') is the cosine (sine) of the angle θ_i' and δ' is a phase responsible for CP violation. We emphasize that these angles are not related to the angles θ_i and phase δ which appear in the Kobayashi-Maskawa[2] matrix K in a completely analogous fashion for three doublets of quarks. For later use, we write the K matrix explicitly as[8]

$$K = \begin{pmatrix} c_1 & s_1c_3 & s_1s_3 \\ -s_1c_2 & c_1c_2c_3+s_2s_3e^{i\delta} & c_1c_2s_3-s_2c_3e^{i\delta} \\ -s_1s_2 & c_1s_2c_3-c_2s_3e^{i\delta} & c_1s_2s_3+c_2c_3e^{i\delta} \end{pmatrix}_{\mathrm{KM}}, \tag{4}$$

where it obviously appears as a natural extension of the Cabibbo-Glashow-Iliopoulos-Maiani 2×2 matrix with δ the phase responsible for CP violation. We append the subscripts Y and KM on the mixing matrices to emphasize the distinction.

There are a number of ways to couple the Higgs doublets to the fermions such that flavor conservation is not violated by neutral-Higgs-boson exchange. Let $\mathcal{P}_i'$, $\mathfrak{N}_i'$, l_i', and ν_i' be the charge $\frac{2}{3}$ quarks, charge $-\frac{1}{3}$ quarks, charged leptons, and neutrinos, respectively, all of them weak-interaction eigenstates. The following cases are typical examples of possible coupling schemes:

(A) ϕ_1 couples to $\mathfrak{N}_{Ri}'$; ϕ_2 couples to $\mathcal{P}_{Ri}'$; ϕ_3 couples to l_{Ri}'. This has discrete symmetry under the transformation: $\phi_2 \to -\phi_2$; $\mathcal{P}_{Ri} \to -\mathcal{P}_{Ri}$; all other fields unchanged. (Similar discrete symmetries exist for ϕ_1 and ϕ_3.) The subscripts R and L denote multiplication with $(1+\gamma_5)/2$ and $(1-\gamma_5)/2$, respectively.

(B) ϕ_1 coupled to $\mathfrak{N}_{Ri}'$, $\mathcal{P}_{Ri}'$, and l_{Ri}'; ϕ_2 and ϕ_3 do not couple to fermions. This has discrete symmetry under the transformation: $\phi_2 \to -\phi_2$ and/or $\phi_3 \to -\phi_3$, all other fields unchanged.

We remark that it is also possible to construct models which allow flavor violation by neutral Higgs bosons for the heavy-quark (b,t) sector. For simplicity, we shall not entertain such possibilities here.

For case (A), the Lagrangian for the Yukawa couplings is given by

$$\begin{aligned}\mathcal{L}_Y = \sum_{i,j=1}^{3} \Big[& \Gamma_{ij}^{(1)}\overline{\mathfrak{N}}_{iR}'(\phi_1^{+*}\mathcal{P}_{jL}' + \phi_1^{0*}\mathfrak{N}_{jL}') \\ & + \Gamma_{ij}^{(2)}\overline{\mathcal{P}}_{iR}'(\phi_2^0\mathcal{P}_{jL}' - \phi_2^+\mathfrak{N}_{jL}') \\ & + \Gamma_{ij}^{(3)}\overline{l}_{iR}'(\phi_3^{+*}\nu_{jL} + \phi_3^{0*}l_{jL}')\Big] + \mathrm{H.c.},\end{aligned} \tag{5}$$

where the $\Gamma_{ij}^{(k)}$ are arbitrary. With spontaneous symmetry breaking, the ϕ_k^0 develop vacuum expectation values λ_k. This yields mass matrices which we diagonalize following the work of Kobayashi and Maskawa. The charged-Higgs-boson Yukawa coupling to the fermion mass eigenstates can be written

$$\mathcal{L}(\phi^\pm)=\overline{\mathcal{P}}_L K M_{\mathfrak{N}} \mathfrak{N}_R \frac{\phi_1^+}{\lambda_1} - \overline{\mathcal{P}}_R M_{\mathcal{P}} K \mathfrak{N}_L \frac{\phi_2^+}{\lambda_2} + \overline{\nu}_L M_l l_R \frac{\phi_3^+}{\lambda_3} + \text{H.c.}, \qquad (6)$$

where

$$\mathfrak{N}=\begin{bmatrix} d \\ s \\ b \end{bmatrix}, \quad \mathcal{P}=\begin{bmatrix} u \\ c \\ t \end{bmatrix}, \quad l=\begin{bmatrix} e \\ \mu \\ \tau \end{bmatrix}, \quad \nu=\begin{bmatrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{bmatrix},$$

and $M_{\mathfrak{N}}, M_{\mathcal{P}}, M_l$ are diagonal mass matrices: $M_{\mathfrak{N}}=\text{diag}(m_d, m_s, m_b)$, $M_{\mathcal{P}}=\text{diag}(m_u, m_c, m_t)$, and $M_l=\text{diag}(m_e, m_\mu, m_\tau)$. K is the Kobayashi-Maskawa matrix, as given in Eq. (4). Here the phases of the fermions are all fixed; however, $\mathcal{L}(\phi^\pm)$ is still invariant under a phase transformation of the Higgs fields. This is because as $\phi_j \to e^{i\Psi_j}\phi_j$, the vacuum expectation value λ_j has a corresponding phase transformation so that ϕ_j/λ_j is invariant. This explains the reason why there is only one physical phase in the Y matrix of Eq. (3).

From the gauge sector, it is straightforward to show

$$\frac{1}{2\sqrt{2}G_F}=\lambda_1^*\lambda_1+\lambda_2^*\lambda_2+\lambda_3^*\lambda_3\,, \qquad (7)$$

while from Eqs. (2), (3), and (7), we have (primes on the angles have been dropped from here on)

$$\begin{aligned} 2^{3/4}G_F^{\frac{1}{2}}\lambda_1 &= (c_1)_Y\,, \\ 2^{3/4}G_F^{\frac{1}{2}}\lambda_2 &= -(s_1c_2)_Y\,, \qquad (8) \\ 2^{3/4}G_F^{\frac{1}{2}}\lambda_3 &= -(s_1s_2)_Y\,. \end{aligned}$$

To write the Yukawa interactions in terms of the physical massive charged Higgs particles, we diagonalize the Higgs bosons using Eq. (2), the Y matrix and Eq. (8). This gives $\mathcal{L}(\phi^\pm)$ in the form

$$\begin{aligned}\mathcal{L}(H^\pm_{1,2}) = {} & 2^{3/4}G_F^{1/2}\overline{\mathcal{P}}_L K M_{\mathfrak{N}}\mathfrak{N}_R\left(\frac{s_1c_3H_1^\pm+s_1s_3H_2^\pm}{c_1}\right)_Y + 2^{3/4}G_F^{1/2}\overline{\mathcal{P}}_R M_{\mathcal{P}} K\mathfrak{N}_L\left[\frac{(c_1c_2c_3+s_2s_3e^{i\delta})H_1^+ + (c_1c_2c_3-s_2c_3e^{i\delta})H_2^+}{s_1c_2}\right]_Y \\ & - 2^{3/4}G_F^{1/2}\overline{\nu}_L M_l l_R\left[\frac{(c_1s_2c_3-c_2s_3e^{i\delta})H_1^+ + (c_1s_2s_3+c_2c_3e^{i\delta})H_2^+}{s_1s_2}\right]_Y + \text{H.c.}, \qquad (9)\end{aligned}$$

and $H_1^\pm$ and $H_2^\pm$ are mass eigenstates with masses m_{H_1} and m_{H_2}, respectively. The charged Goldstone boson has been absorbed into the W gauge field and hence is not displayed in Eq. (9). All the c_i, s_i, and δ shown explicitly in Eq. (9) belong to the Y matrix.

We observe that, as written here, the CP-violating effects are all contained in the couplings of physical particles (i.e., quark–massive-gauge-field couplings, quark–charged-Higgs-boson couplings, Higgs-boson self-couplings, and Higgs-boson–massive-gauge-field couplings). For the purpose of applications this way of casting the CP-violating effects is superior to the original approach, where CP invariance is violated by the Higgs propagators.[5] As opposed to the Kobayashi-Maskawa model, CP violation now occurs in the lepton sector as well as in the quark sector. Hence the CP violation generated by the Higgs mechanism can, in principle, be distinguished from the KM model by experiment. The neutral Higgs sector can be treated in a similar fashion. Extension to n Higgs doublets and m quark (and lepton) doublets is straightforward.

Finally we remark for case (B), where ϕ_2 and ϕ_3 do not couple to fermions, the charged-Higgs-particle coupling to fermions has the form

$$\mathcal{L}(H^\pm)=2^{3/4}G_F^{1/2}\left[\overline{\mathcal{P}}_L K M_{\mathfrak{N}}\mathfrak{N}_R - \overline{\mathcal{P}}_R M_{\mathcal{P}} K \mathfrak{N}_L + \overline{\nu}_L M_l l_R\right] \times\left(\frac{s_1c_3H_1^+ + s_1s_3H_2^+}{c_1}\right)_Y. \qquad (10)$$

The CP-violating effects from the Higgs sector now remain in the Higgs-boson self-coupling sector.

III. CHARGED-HIGGS-BOSON PHENOMENOLOGY

We now take up the subject of charged-Higgs-boson phenomenology and restrict our attention to just one aspect of this problem, namely, ways to carry out a search for charged Higgs bosons in e^+e^- annihilation within the energy range accessible to forthcoming experiments. By definition, $m_{H_2} \geq m_{H_1}$, and for simplicity, we assume that H_2 has a much larger mass than H_1 and consider just the light one, $H_1^\pm$, which we relabel $H^\pm$. In the presence of new heavy quarks and leptons, the charged-Higgs-boson search can be carried out via their decays as discussed in the literature. Here we shall concentrate on other search techniques.

The couplings of $H^\pm$ to the quarks and leptons can be read off immediately from Eqs. (9) or (10).

While bounds have been placed on the KM angles ϕ_1, ϕ_2, ϕ_3 and δ, the Y-matrix angles for the Higgs rotation remain unrestricted at this point. In principle, experimental bounds on the Higgs-boson masses and Y-matrix angles can be obtained, but we expect only very loose constraints and do not derive them here. In any event, it is clear by virtue of the τ mass that among the leptons, τ and ν_τ couple most strongly to the Higgs boson $H^\pm$ for models with just three generations of lepton-doublets. This suggests that we define a parameter r

$$r=\frac{(g_{q\bar{q}H})_{\text{max effective}}}{g_{\nu_\tau \tau H}}, \tag{11}$$

which corresponds to the ratio of the maximum effective strength of the quark-Higgs-boson coupling to the strength of the $\nu_\tau-\tau-H$ coupling. From Eq. (9) it is clear that r can be much larger than, roughly equal to, or much smaller than unity. We shall, therefore, consider the three cases where $r\ll 1$, $r\sim 1$, and $r\gg 1$.

In Table I we present the most favorable ways to search for charged Higgs bosons in e^+e^- annihilation for different possible mass ranges and for each of the three r cases. Only the mass range $m_H<m_\tau$ seems to have been excluded by the present data since the branching ratios for τ decays do not allow decays into charged Higgs bosons. We now discuss each of the other entries in the table in turn. For simplicity, we shall ignore the CP-violation phase, since CP-violating effects do not provide easy direct signals for the search of charged Higgs bosons.

A. $m_H \sim m_{D,F}$

If m_{H^+} is comparable to the charmed-meson masses and $r\lesssim 1$, the D^+ and F^+ mesons will exhibit enhanced decay rates into the $\tau\nu_\tau$ channel. Using the Higgs-boson coupling for case (A), we have, in the approximation where $m_c \gg m_s, m_d$ and where strong-interaction corrections are ignored,

$$\frac{\Gamma(D^+\to H^+\to\tau^+\nu_\tau)}{\Gamma(D^+\to W^+\to\tau^+\nu_\tau)}\simeq\frac{m_c^2 m_\tau^2}{(m_H^2-m_{D^+}^2)^2}\times\left[\frac{(c_1c_2c_3+s_2s_3)(c_1s_2c_3-c_2s_3)}{s_1^2c_2s_2}\right]^2_Y \simeq\frac{P}{(m_H^2-m_{D^+}^2)^2} \tag{12}$$

and

$$\frac{(F^+\to H^+\to\tau^+\nu_\tau)}{(F^+\to W^+\to\tau^+\nu_\tau)}\simeq\frac{P}{(m_H^2-m_{F^+}^2)^2}\ . \tag{13}$$

TABLE I. Signatures for charged Higgs bosons in e^+e^- annihilation for various mass ranges and values of r defined in Eq. (11).

	$r\ll 1$	$r\sim 1$	$r\gg 1$
$m_H<m_\tau$	Ruled out by $\tau^+\to H^++\bar{\nu}_\tau$		Ruled out by $\psi\to H^++H^-+q+\bar{q}$ (see Ref. 10)
$m_H\sim m_{D,F}$	$D^+,F^+\to H^+\to\tau\nu_\tau$ enhancement		$e^+e^-\to(b\bar{b})\to H^\pm+X$ $e^+e^-\to(b\bar{b})\to H^++H^-+q\bar{q}$ (see ref. 10)
$m_F<m_H<m_B$	$\Delta R\lesssim 0.25$ in $\mu e,\mu X$, eX channels $e^+e^-\to H^+H^-$ $H^-\to\tau\bar{\nu}_\tau$, $\tau\to e^-,\mu^-,X^-+\nu_\tau$ $H^+\to\tau^+\nu_\tau$, $\tau^+\to e^+,\mu^+,X^++\bar{\nu}_\tau$	$\Delta R\lesssim 0.1$ $e^+e^-\to H^+H^-$ $H^-\to$ hadronic jet $H^+\to\tau^+\nu_\tau$, $\tau^+\to e^+,\mu^+,X^++\bar{\nu}_\tau$	Rise in ΔR in the multikaon channel (see Ref. 11)
$m_H\sim m_B$	$B^+\to H^+\to\tau\nu_\tau$ enhancement		Rise in ΔR in the multikaon channel (see Ref. 11)
$m_H\gg m_B$	$\Delta R\lesssim 0.25$ in μe, μX, eX channels $e^+e^-\to H^+H^-$ $H^-\to\tau^-\bar{\nu}_\tau$, $\tau^-\to e^-,\mu^-,X^-+\nu_\tau$ $H^+\to\tau^+\nu_\tau$, $\tau^+\to e^+,\mu^+,X^++\bar{\nu}_\tau$	$e^+e^-\to H^+H^-$ $H^-\to$ hadronic jet $H^+\to\tau^+\nu_\tau$, $\tau^+\to e^+,\mu^+,X^++\bar{\nu}_\tau$	Rise in ΔR in the multikaon channel (see Ref. 11)

The explicit widths for the D^+,F^+ decays via the H^+ depend on the values of the Y matrix angles and the mass of H^+. In general, we expect an enhancement of the rate of the $\tau^+\nu_\tau$ decay mode. The branching ratio of the $\tau^+\nu_\tau$ decay (via W^+) of the Cabibbo-favored F^+ case is ~3%, while that for the Cabibbo-suppressed D^+ case is 0.02%.[9] If H^+ roughly gives the same enhancement factor to the F^+,D^+ decays into $\tau^+\nu_\tau$, say, a factor of 20, then the D^+ decay branching ratio to $\tau^+\nu_\tau$ is still only 4% while the F^+ decays predominantly to $\tau^+\nu_\tau$. This would have the distinguishing feature of making it difficult to observe the F^+ decay.

For the case where $r \gg 1$ and $m_H \sim m_{D,F}$, it has been suggested[10] that charged Higgs bosons could be pair produced via the decay of a heavy quarkonium bound state (e.g., Υ) via the mode $(Q\overline{Q}) \to H^+H^- + q\bar{q}$, which has the distinct feature of having a large acoplanarity. For $r \gtrsim 1$, we expect to see events from the quarkonium decay which have both leptons and hadrons in the final state. The decay mode $(Q\overline{Q} \to H^+$ + hadrons may also be a suitable place for the search of charged Higgs bosons. In fact, $(Q\overline{Q}) \to H^\pm$ + anything can be an important decay mode of the quarkonium state.

B. $m_F < m_H < m_B$

For this mass range of 2–5 GeV, charged Higgs bosons could be pair produced in e^+e^- annihilation at SPEAR, DORIS, and CESR. Since the Higgs particles are scalars, the reaction $e^+e^- \to H^+H^-$ would add at most one quarter of a unit to R. With $r \gg 1$, it has already been suggested that one try to observe a change in R in the multikaon channels.[11]

For $r \ll 1$, one can look for the process

$e^+e^- \to H^+H^-$

$H^- \to \tau^-\bar{\nu}_\tau \to$ one charged track

$H^+ \to \tau^+\nu_\tau \to$ one charged track. (14)

This gives an increase in R roughly of the size $\Delta R \lesssim 0.2$ in the two-charged-particle final state. In Fig. 1, we present distributions for (a) the energy of the μ resulting from the $H \to \tau \to \mu$ chain, (b) the angle of the μ relative to the beam direction, (c) the angle of the μ relative to the H production direction, and finally (d) the opening angle between the muon and electron produced in the decay chains of the two Higgs bosons. We have chosen $m_H = 2.5$ GeV and $E_{\text{beam}} = 3.5$ GeV for purposes of illustration. The lepton distributions resulting from the direct pair production of τ leptons are also given for comparison. We observe that

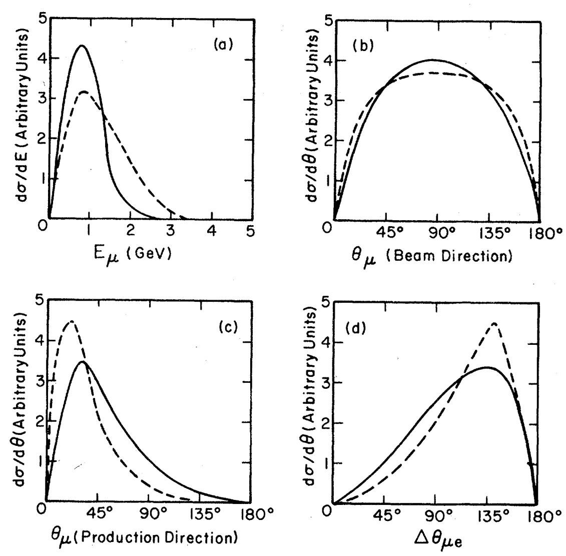

FIG. 1. Muon (a) energy, (b) angle with respect to beam direction, (c) angle with respect to production direction, and (d) opening angle between muon and electron for 2.5-GeV charged-Higgs-boson pair production (solid curves) and 1.78-GeV τ pair production (dashed curves) in e^+e^- annihilation with $\sqrt{s} = 7$ GeV. The absolute values of muon production from H^+ and τ^+ decays in e^+e^- annihilation are obtained by multiplying the curves by the appropriate cross section time branching ratios after normalizing the integrated distributions to unity.

the signals are not significantly different for the Higgs-boson-pair and τ pair processes, so it will be difficult indeed to identify a small Higgs-boson signal under a large τ background. It may be useful to look for a rise in R in this channel only in certain kinematic regions so as to minimize the direct $\tau^+\tau^-$ production background.

C. $m_H \sim m_B$

For Higgs-boson masses close to the B-meson masses, one would expect to obtain a Higgs propagator enhancement of the form $m_\tau^2 m_b^2/(m_H^2 - m_B^2)$ in the decay mode $B^+ \to H^+ \to \tau^+\nu_\tau$ similar to that in (12) and (13) for the F and D decays.

D. $m_H \gg m_B$

If the Higgs-boson mass is much larger than the B-meson masses, one can use the same search techniques discussed in (2) above, now in the energy range spanned by PETRA and PEP. In the case where $r \lesssim 1$, we have illustrated the expected muon distributions in Fig. 2 for $E_{\text{beam}} = 20$ GeV and $m_H = 5$, 10, and 15 GeV. The results from the Higgs-boson-pair chain decays of Eq. (14) are now significantly different from the τ-pair background, especially θ_μ relative to the H production direction

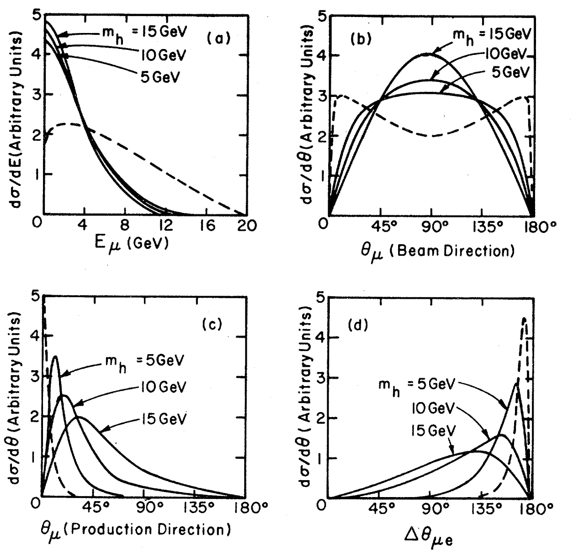

FIG. 2. Distributions as in Fig. 1 for a charged-Higgs-boson mass of 5, 10, and 15 GeV and $\sqrt{s} = 40$ GeV with the dashed curves again referring to τ pair production.

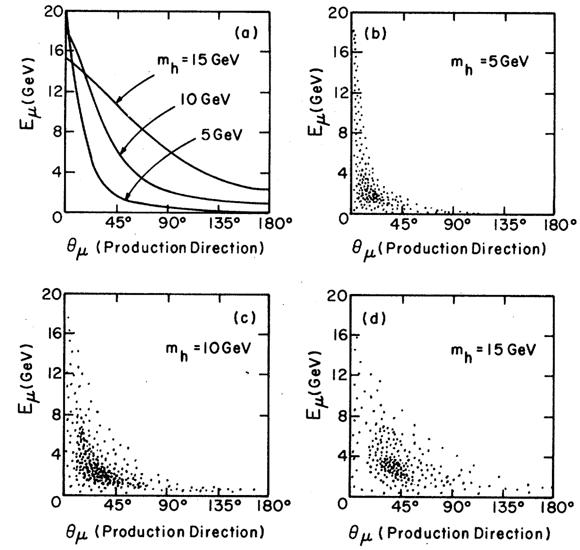

FIG. 3. Plots of E_μ vs θ_μ with respect to production direction plots with (a) showing the kinematical boundaries and (b)–(d) scatter plots with charged-Higgs-boson masses of 5, 10, and 15 GeV, respectively, and $\sqrt{s} = 40$ GeV.

and $\Delta\theta_{\mu e}$. However, the Higgs-boson signals are suppressed by a factor of $\gtrsim 8$ relative to the τ signals owing to factors of 0.25 for production, $\lesssim 2(0.5)^2$ for the $H \to \tau$ branching ratio, and less than unity for the branching ratio of the τ into the desired detection channel.

This suggests that the most appropriate means of separating a Higgs-boson-pair signal from a τ-pair signal is to search for μ-hadron or e-hadron events in which the invariant mass of the hadron jet (coming from one H decay) is larger than the τ mass:

$$e^+e^- \to H^+H^-,\quad H^- \to \tau^-\bar{\nu}_\tau,\quad H^+ \to \text{hadrons}. \tag{15}$$

The signal should then be very clean and can be compared, in particular, with the curves presented in Figs. 2(a) and 2(c). In Fig. 3(a) we show the boundary curves for E_μ vs. θ_μ with respect to the production direction (defined to be opposite the hadron jet direction) and scatter plots in these variables in Figs. 3(b), 3(c), and 3(d) for $m_H = 5$, 10, and 15 GeV, respectively. Similar curves and scatter plots will obtain for any E_{beam} and m_H in the region considered such that $m_H/E_{\text{beam}} = 0.25$, 0.50, and 0.75. Typically, both E_μ and θ_μ can be relatively large for Higgs-boson pair production near threshold.

E. $m_H > m_\tau$, $r \gtrsim 1$

In general, charged Higgs bosons would enhance the decay modes that involve heavy quarks. For example, Kane[12] has suggested that the existence of charged Higgs bosons can, in general, enhance the ratio $\Gamma(D^0 \to K^+K^-)/\Gamma(D^0 \to \pi^+\pi^-)$. Also we expect that, for the B system, the branching ratios of the decay modes that involve the charm quark, e.g., $\Gamma(B^- \to \psi K^-)$, $\Gamma(B^0 \to \psi K^-\pi^+)$, $\Gamma(B^0 \to D\bar{D}K^-\pi^+)$, would be enhanced with respect to that of the decay modes which do not involve the charm quark. (Such an enhancement effect should be even more prominent in the Υ-meson decays into final states which include a $b\bar{b}$ pair.) As a typical example, the branching ratio of the decay mode $B^0 \to \psi K^-\pi^+$ should be enhanced by a factor of $[1 + Am_b^2m_c^2/(m_H^2 - M^2)^2]$, where A is a constant of order 1 and M is a typical invariant mass of $\bar{c}s$, $M \sim 2$ GeV. A more precise estimate is beyond the scope of this paper.

To summarize, we have extended the various signals for charged-Higgs-boson searches suggested in the literature. Although the arguments for the existence of charged Higgs bosons are in no way compelling, we believe it is important to experimentally confirm or disprove the existence of charged Higgs bosons in given mass ranges. As is clear from the table, if charged Higgs bosons are discovered, the particular mode where it is found will also provide useful information concerning its coupling properties. Details of the Higgs-boson-induced CP-violation phenomenology will be discussed elsewhere.

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