Higgs-boson associated production in e^+e^- annihilation below the Z

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The associated production of a scalar (H) and pseudoscalar (P) Higgs boson at the e^+e^- collider TRISTAN at KEK is discussed within the framework of minimal supersymmetry. We first discuss constraints on the model from the nonobservation of radiative decays of the Υ . We then show that a substantially larger region of parameter space can be probed already at $\sqrt{s} = 60$ GeV, where a total HP cross section in excess of 1 pb is possible. We discuss the clean $\tau^+\tau^-$ pair signal in some detail and show how its opening angle can be used to reconstruct the mass of the parent Higgs boson. Some remarks on purely hadronic decay final states are also presented.

I. INTRODUCTION

One of the most important tasks of contemporary high-energy physics is to gain insight into the mechanism that gives masses to W and Z bosons. One appealingly simple explanation postulates the existence of spinless Higgs fields¹ with a nonvanishing vacuum expectation value (VEV) which spontaneously breaks the local $SU(2) \times U(1)_Y$ gauge invariance governing the interactions of W and Z bosons with matter fermions and with each other. In the simplest possible case, the standard model (SM), only one complex $SU(2)$ doublet of Higgs bosons is introduced; three of these four degrees of freedom are absorbed by the longitudinal gauge bosons, so that only one physical Higgs field, a neutral scalar, remains.

Unfortunately, the SM is technically unnatural:² Quantum corrections tend to destabilize the scale at which electroweak symmetry breaking occurs, driving it up to some ultrahigh scale,³ such as the scale M_X at which the hypothetical grand unified gauge symmetry is broken, or even the Planck scale M_P . The best understood solution of this problem relies on the introduction of supersymmetry; quantum corrections involving the new superpartners cancel⁴ the bad behavior of corrections arising from SM fields. This formally elegant solution, however, necessitates the introduction of a multitude of new particles, none of which has yet been found; not only does every SM field get a superpartner, the Higgs sector itself, source of all the problems of the SM, has to be enlarged⁵ to contain at least two $SU(2)$ doublets of Higgs superfields.

This theoretical complication of the Higgs sector, however, might be a blessing from the experimental point of view, since it leads to production mechanisms for Higgs bosons that are absent in the SM. In this paper we focus on one of these new mechanisms, the associated production of a neutral Higgs scalar H with a pseudoscalar P at the e^+e^- collider TRISTAN at KEK. We will work in the framework of minimal supersymmetry,⁶ where the whole Higgs sector can be parametrized in terms of two

unknown parameters, but our results should remain qualitatively correct in supersymmetric models⁷ with a more complicated Higgs sector, or even in nonsupersymmetric models⁸ with at least two Higgs doublets.

In fact, the strongest reason in adopting the minimal Higgs sector in the SM (one doublet) is the simplicity of the model. It is almost the *only* possible Higgs sector which does *not* allow for the associated HP production, the only neutral pseudoscalar being absorbed by the Z boson. In any other model a HP pair can be produced in e^+e^- annihilation via Z exchange. The production characteristics are rather model independent, since the production mechanism is determined by the gauge structure up to a mixing factor which will be discussed later. Therefore, most conclusions of the paper are widely applicable at least qualitatively.

The remainder of this paper is organized as follows. In Sec. II we briefly review the model and discuss constraints on its parameters from the negative search for the Wilczek decay of the Υ vector meson into a Higgs boson and a photon. In Sec. III we present the total cross section for HP production, emphasizing the big gain of going from PETRA (at DESY) to TRISTAN energies, where cross sections in excess of 1 pb are possible. In Sec. IV we discuss signals and event reconstruction. We concentrate on the case where at least one of the Higgs bosons decays into τ leptons, since this signal should be easily recognizable and almost free of (physics) background. We also briefly discuss the case where both H and P decay into quarks. We show that in either case the opening angle between the decay products of one Higgs boson is strongly peaked at its minimum possible value, allowing for a determination of the mass of this boson. Finally in Sec. V we summarize our findings.

II. EXISTING CONSTRAINTS ON MODEL PARAMETERS

As stated in the Introduction, we work within the framework of minimal supersymmetry, where the Higgs

sector consists of just two SU(2) doublets. In this model the masses and couplings of all five physical degrees of freedom (one charged scalar $H^\pm$, one neutral pseudoscalar P , and two neutral scalars H^0, H'^0) can be computed⁹ as functions of just two unknown parameters. Since the charged and heavier neutral Higgs bosons are always heavier than W and Z , respectively, we are here only interested in the pseudoscalar and the lighter neutral Higgs boson, which we will simply call H from now on. As the two independent parameters we can then choose the ratio ω of the VEV's of the two neutral Higgs fields plus either m_P or m_H . These three quantities are related by⁹

$$M_H^2 = \frac{1}{2} \left\{ m_P^2 + m_Z^2 - \left[(m_P^2 + m_Z^2)^2 - 4m_P^2 m_Z^2 \left(\frac{\omega^2 - 1}{\omega^2 + 1} \right)^2 \right]^{1/2} \right\}, \quad (2.1)$$

which immediately implies

$$m_H < \left| \frac{\omega^2 - 1}{\omega^2 + 1} \right| \min(m_Z, m_P). \quad (2.2)$$

A third choice of independent parameters, m_P and m_H , is not sufficient to fix the model in principle, since Eq. (2.1) is invariant under $\omega \rightarrow 1/\omega$. However, one usually assumes the VEV giving rise to the top-quark mass to be bigger than the one giving the mass of the bottom quark,¹⁰ which means $\omega > 1$; under this assumption m_H and m_P do fix the model uniquely.

The most restrictive existing experimental bound¹¹ on the Higgs sector stems from the nonobservation of the Wilczek decay¹² of the vector ($b\bar{b}$) state, $\Upsilon \rightarrow A + \gamma$ ($A = H, P$). Unfortunately, the interpretation of this experimental null result is not straightforward. First, one has to assume that this decay can be adequately described within the framework of nonrelativistic Van Royen–Weisskopf formalism based on the static-quark approximation;¹³ it has been shown¹⁴ to break down if m_A comes close to m_Υ for the case $A = H$ (and to a much lesser extent for $A = P$). Second, even within the static-quark formalism, QCD corrections to the decay width turn out^{15,16} to be large and negative (especially in the scalar case, $A = H$), and even make the width negative for small values of the photon energy.

In leading order, the relevant decay widths are given in our case:

$$\Gamma_0(\Upsilon \rightarrow P + \gamma) = \frac{\alpha_{\text{em}} G_F}{9\sqrt{2}\pi} |R_S(0)|^2 \omega^2 \left[1 - \frac{m_P^2}{m_\Upsilon^2} \right], \quad (2.3)$$

$$\Gamma_0(\Upsilon \rightarrow H + \gamma) = \frac{\alpha_{\text{em}} G_F}{9\sqrt{2}\pi} |R_S(0)|^2 \left[\frac{\sin\alpha}{\cos\beta} \right]^2 \left[1 - \frac{m_H^2}{m_\Upsilon^2} \right], \quad (2.4)$$

where $R_S(0)$ is the Υ wave function at the origin, $\tan\beta \equiv \omega$, and α is an angle describing the mixing between the neutral Higgs scalars which can be computed from⁹

$$\tan 2\alpha = \frac{m_P^2 + m_Z^2}{m_P^2 - m_Z^2} \tan 2\beta, \quad (2.5)$$

with the constraint $-\pi/4 \leq \alpha \leq 0$ (if $\omega < 1$) or $-\pi/2 \leq \alpha \leq -\pi/4$ (if $\omega > 1$) in the (for us) relevant case $m_P < m_Z$. The one-loop QCD corrections can be written as

$$\Gamma_1(\Upsilon \rightarrow A + \gamma) = \Gamma_0(\Upsilon \rightarrow A + \gamma) \times \left[1 - \frac{4}{3\pi} \alpha_s (m_\Upsilon^2 - m_A^2) a_A \left(\frac{m_A^2}{m_\Upsilon^2} \right) \right], \quad (2.6)$$

where the $a_A(x)$ are complicated functions,^{15,16} which can, however, be approximated accurately by

$$a_H(x) \simeq 5.74\sqrt{1-x} + \frac{4.19}{\sqrt{1-x}}, \quad (2.7)$$

$$a_P(x) \simeq 2.00 + 4.62(1-x)^{0.28}. \quad (2.8)$$

The value for $|R_S(0)|^2$ is usually obtained by equating the measured $\Upsilon \rightarrow e^+e^-$ decay width¹⁷ with the prediction of the static-quark model, where for consistency one again has to include one-loop α_s corrections.¹⁸ This gives

$$|R_S(0)|^2 = \Gamma_{\text{expt}}(\Upsilon \rightarrow e^+e^-) \times \frac{9m_\Upsilon^2}{4\alpha_{\text{em}}^2} \times \left\{ \left[1 - \frac{16\alpha_s(m_\Upsilon^2)}{3\pi} \right]^{-1} \left[1 + \frac{16\alpha_s(m_\Upsilon^2)}{3\pi} \right] \right\}. \quad (2.9)$$

These two expressions are equivalent up to $O(\alpha_s^2)$; however, when plugged into Eqs. (2.4) and (2.6), they lead to quite different results. This difference is usually taken to be the measure of the size of (yet unknown) two-loop corrections. To be even more conservative, we use the upper form of Eq. (2.9) with a large $\Lambda_{\overline{\text{MS}}} = 300$ MeV, and the lower with a small $\Lambda_{\overline{\text{MS}}} = 150$ MeV ($\overline{\text{MS}}$ denotes the modified minimal-subtraction scheme). We finally parametrize the experimental bound as

$$\Gamma(\Upsilon \rightarrow A + \gamma) < 0.52 \times 10^{-2} \text{ keV} \left[0.28 + 0.14 \frac{m_A}{1 \text{ GeV}} \right], \quad (2.10)$$

for $m_A \lesssim 8.5$ GeV, where we have used¹⁷ $\Gamma_{\text{tot}}(\Upsilon) = 52$ keV.

The excluded region¹⁹ of the m_H - ω plane is shown in Fig. 1. The region between the dotted lines is unphysical; see Eq. (2.2). The solid and dashed curves represent the upper and lower choice in Eq. (2.9), respectively. We see that the uncertainty due to α_s^2 terms is indeed sizable in the scalar case, but not so big that one would lose all confidence in the results of Fig. 1. Furthermore, the bound from $\Upsilon \rightarrow P + \gamma$ is much less affected by this uncertainty. Note, however, that the bound has a

FIG. 1. Excluded regions in the (m_H, ω) plane. The region between the dotted lines is excluded by Eq. (2.1). The region to the left and above the curves labeled P and H is excluded by the nonobservation of $\Upsilon \rightarrow P + \gamma$ and $\Upsilon \rightarrow H + \gamma$, respectively, where the dashed and solid lines correspond to different treatments of QCD corrections as discussed in the text below Eq. (2.9). The area within the parabola-shaped curves at $m_H \approx 1$ GeV, $\omega \approx 0.7$ is also excluded from $\Upsilon \rightarrow P + \gamma$.

confidence level of only 90%. Furthermore, all Higgs-boson masses and mixings have been computed at the tree level; this is certainly inadequate as $m_{H,P} \rightarrow 0$. Several groups²⁰ have studied the effect of finite radiative corrections in the manner of Coleman and Weinberg²¹ to m_H and the scalar potential, usually focusing onto the extreme case where $m_H = 0$ at the tree level. One finds that the corrections to m_H are somewhat smaller than in the SM (Ref. 22) (with a light top quark), and can even be zero, but are in general of the order of 3–6 GeV. No computation of corrections to m_P exist, but one might expect them to be similar in magnitude. Figure 1 is thus very unreliable²³ in the region $m_H \lesssim 5$ GeV. In view of all these uncertainties we may, therefore, conclude that it might be premature to really rule out any region of parameter space of our model. Of course, this leaves the door wide open for TRISTAN to investigate the model.

III. TOTAL CROSS SECTION FOR $e^+e^- \rightarrow HP$

Since H and P have no electric charge, they can only be produced by exchanging a Z boson in the s channel. Furthermore, HH and PP production is forbidden by Bose symmetry. Associated HP production is, however, possible;²⁴ the total cross section is given by

$$\sigma(e^+e^- \rightarrow HP) = \frac{2\pi\alpha_{\text{em}}^2(v_e^2 + a_e^2)}{3\sin^4\theta_W \cos^4\theta_W} \times \cos^2(\alpha - \beta) \frac{k^3}{\sqrt{s}[(s - m_Z^2)^2 + m_Z^2\Gamma_Z^2]}, \quad (3.1)$$

where $v_e = -\frac{1}{4} + \sin^2\theta_W$ and $a_e = -\frac{1}{4}$ describe the vector and axial-vector coupling of electrons to Z bosons, and

the three-momentum k of H (or P) in the center-of-mass system is given by

$$k^2 = \frac{1}{4s}[(s - m_H^2 - m_P^2)^2 - 4m_H^2m_P^2]. \quad (3.2)$$

Finally, the suppression factor $\cos^2(\alpha - \beta)$ is given by

$$\cos^2(\alpha - \beta) = \frac{m_H^2}{m_P^2} \frac{m_Z^2 - m_H^2}{m_Z^2 + m_P^2 - 2m_H^2}. \quad (3.3)$$

This factor describes the “matching” of H and P , and vanishes if $m_H \rightarrow 0$ for fixed m_P (which corresponds to $\omega \rightarrow 1$), and approaches unity as $m_H \rightarrow m_P$, i.e., $\omega \rightarrow 0$ or ∞ (recall that $m_H \leq m_P$).

In Fig. 2 we show the energy dependence of the total cross section for HP production in the most favorable case $m_H = m_P$. Notice the large increase of the cross section as going from PETRA energies ($\sqrt{s} \lesssim 46$ GeV) to TRISTAN ($\sqrt{s} \approx 60$ GeV). Even if there is no big phase-space suppression, $m_H = m_P = 10$ GeV, the cross section rises by a factor of 3.5; for $m_H = m_P = 20$ GeV, one gains almost a factor of 10. At the same time, many potential backgrounds, which involve the exchange of a photon in the s channel, continue to decrease. Notice also that a total HP cross section in excess of one picobarn is possible already at $\sqrt{s} = 60$ GeV.

The differential cross section is proportional to $\sin^2\theta$, where θ is the angle between one of the Higgs bosons and the beam axis in the c.m. frame. This implies relatively few loss of events in the forward/backward region as compared to fermion-pair production.

In Fig. 3, we show lines of constant total cross section for HP production at $\sqrt{s} = 60$ GeV. We see that the cross section falls off quite rapidly as one moves away from the line $m_H = m_P$. We expect this feature to survive even for nonminimal models with additional SU(2)-singlet Higgs superfields. In this case, the HP cross section might, however, be small even at $m_H = m_P$, if H and P are dominantly SU(2) $\times$ U(1) singlets. On the other hand, already with an integrated luminosity of only 10 pb^{-1} one can probe quite a large region of parameter space of

FIG. 2. Total cross section for $e^+e^- \rightarrow Z^* \rightarrow H + P$ as a function of the center-of-mass energy $\sqrt{s}$. It has been assumed that $m_H = m_P$, which maximizes the ZHP coupling.

FIG. 3. Lines of constant $\sigma(e^+e^- \rightarrow HP)$ at $\sqrt{s}=60$ GeV in the (m_H, m_P) plane. All cross sections are in pb. The region above the straight line $m_H = m_P$ is unphysical in minimal supersymmetry, where one always has $m_H \leq m_P$.

the minimal model, at least in principle. With 50–100 pb, one is sensitive to almost the whole kinematically accessible region, unless $m_H^2 \ll m_P^2$, i.e., $\omega \simeq 1$. A further major increase of the accessible part of parameter space will then only be possible at higher energies.²⁵ In Ref. 26 it has been shown that the forthcoming e^+e^- colliders at the Z^0 energy, LEP (at CERN) and SLC (at SLAC), can test the whole region $m_H \lesssim 40$ GeV by either HP production or the Bjorken process²⁷ $Z \rightarrow Z^*H$. However, our point here is that one may not have to wait for LEP/SLC. Whether TRISTAN can actually participate in the hunt for Higgs bosons depends, of course, on the signals of HP production, which we discuss in the following section.

IV. DECAY MODES AND EVENT SIGNATURE

As is well known, the SM Higgs boson always tends to decay into the heaviest possible particles. This trend still persists in the somewhat more complicated case of minimal supersymmetry, but now one has to include effects due to the mixing of the original two Higgs doublets. For the mass range of interest to us, the only possible tree-level decay modes²⁸ of H and P are into quark-antiquark or lepton-antilepton pairs. The relevant couplings can be written in the general form

$$g_{Af\bar{f}} = \frac{gm_f}{2m_W} C_{Af\bar{f}}, \quad (4.1)$$

where $A=H$ or P , and g is the SU(2) gauge coupling: $g^2/4m_W^2 = \sqrt{2}G_F$. The factor $C_{Af\bar{f}}$ originates from the mixing of the two Higgs doublets and is given by⁹

$$C_{Pu\bar{u}} = 1/\omega, \quad C_{Pd\bar{d}} = C_{Pl\bar{l}} = \omega, \quad (4.2)$$

$$C_{Hu\bar{u}} = \frac{\cos\alpha}{\sin\beta}, \quad C_{Hd\bar{d}} = C_{Hl\bar{l}} = \frac{\sin\alpha}{\cos\beta},$$

where u, d, l stand for any up-type quark, down-type quark, and charged lepton, respectively, and α is given by Eq. (2.5). We thus see that the couplings of both H and P to down-type quarks and charged leptons are enhanced (suppressed) for $\omega > 1$ ($\omega < 1$); indeed one finds that $C_{Hf\bar{f}}$

approaches $C_{Pf\bar{f}}$ as $\omega \rightarrow 0$ or ∞ . The $A \rightarrow f\bar{f}$ decay widths are then given by

$$\Gamma(A \rightarrow f\bar{f}) = N_c \frac{g_{Af\bar{f}}^2}{8\pi} m_A \left[1 - \frac{4m_f^2}{m_A^2} \right]^{l_A}, \quad (4.3)$$

where $N_c = 3$ (1) for quarks (leptons), $l_H = \frac{3}{2}$, and $l_P = \frac{1}{2}$.

Most relevant for Higgs bosons with masses between a few GeV and about 30 GeV, which is the region accessible at TRISTAN, are decays into $c\bar{c}$, $b\bar{b}$, and $\tau^+\tau^-$ pairs. Of these three modes, the decay into $\tau^+\tau^-$ offers the cleanest, most easily recognizable signal. In an ideal case, both H and P would decay into τ pairs. In this case, the efficiency for detecting a τ pair in one hemisphere, with the other hemisphere being consistent with a τ pair also, should be fairly high. Note also that because of the $\sin^2\theta$ angular distribution, the efficiency to detect all the decay products of H and P will be rather high. However, the probability that both H and P decay into τ 's becomes very small if P is heavier than 12 GeV or so. We, therefore, only require H or P to decay into a τ pair and do not impose any constraints on the second Higgs boson.

In Fig. 4 we show lines of constant $\sigma(e^+e^- \rightarrow HP \rightarrow \tau^+\tau^- X)$ at $\sqrt{s}=60$ GeV, where we have assumed $\omega > 1$. The most prominent feature of these curves is the spike towards large m_P at $m_H \simeq 10$ GeV. For larger m_H , the $H \rightarrow b\bar{b}$ decay becomes possible and quickly dominates $H \rightarrow \tau^+\tau^-$. The resulting decrease of the branching ratio into τ pairs can, however, be balanced by going closer to the line $m_H = m_P$, where the mixing factor is maximal and the total HP cross section is largest (see Fig. 3). We should also mention that our results are not very reliable if m_H or m_P is close to 10 GeV, where threshold effects are important: Just below threshold, the Higgs boson can mix²⁹ with $b\bar{b}$ bound states, while just above threshold the consideration of exclusive decay modes would be necessary to obtain a reliable rate. We also did

FIG. 4. Lines of constant $\sigma(e^+e^- \rightarrow HP \rightarrow \tau^+\tau^- X)$ at $\sqrt{s}=60$ GeV in the (m_H, m_P) plane. All cross sections are in pb. This figure is only valid for $\omega > 1$; for $\omega < 1$, the branching ratios for H or P to $\tau^+\tau^-$ are much smaller, so that the cross section for the τ signal is always below 0.05 pb.

not include QCD corrections, which can be expected to be maximal close to threshold.

Nevertheless, we can conclude from Figs. 3 and 4 that if $\omega > 1$ a sizable fraction of parameter space that is in principle accessible at TRISTAN can indeed be tested using the τ pair signal. The background to this signal should be fairly small, if we require that the system X in the hemisphere opposite to the τ pair is not just a photon (to exclude radiative $\tau^+\tau^-$ production). In this case all backgrounds are of order α_{em}^4 , and can be further reduced if one requires X to be consistent with being a $c\bar{c}$ or $b\bar{b}$ pair.

One disadvantage of the τ -pair signal is that the invariant mass of the pair cannot be measured directly, since some energy-momentum is always lost in the form of neutrinos. Note, however, that the τ leptons are usually rather fast: Their energy spectrum in the laboratory frame is flat, with end points

$$E_{\pm} = \frac{1}{2}m_A \left[1 + \frac{k^2}{m_A^2} \right]^{1/2} \pm \frac{k}{m_A} p^*, \quad (4.4)$$

where $A = H$ or P is the Higgs boson that decays into τ pairs, k is given by Eq. (3.2), and p^* is the three-

momentum of the τ in the rest frame of A :

$$p^{*2} = \frac{1}{4}m_A^2 - m_{\tau}^2. \quad (4.5)$$

The first term in Eq. (4.4) gives the average τ energy $\bar{E}$. For our purposes, the average energy of the slower of the two τ leptons is perhaps more relevant; it is given by $\bar{E}_{\text{slow}} = \frac{1}{2}(E_- + \bar{E})$ and lies between 7 and 10 GeV as long as $m_P < 30$ GeV. This means that 50% of the slower τ 's will have a γ factor of 4 or more. The only configuration where all τ leptons are likely to be slow comes from the production of a light H and a heavy P with $H \rightarrow \tau^+\tau^-$; meanwhile, the cross section is suppressed in this case both by phase-space constraints [the k^3 factor in Eq. (3.1)], and by the mixing factor in the ZHP coupling $\cos(\alpha - \beta)$, Eq. (3.3). This configuration is thus unlikely to be observable at TRISTAN (see Fig. 4). Since the τ leptons tend to be fast, all their decay products will go more or less into the same direction as the τ itself. A fairly precise measurement of the opening angle of the τ pair should thus be possible. This angle is large for a heavy, slow-moving Higgs boson, and small if the Higgs boson that decays into the τ pair is light and fast. It is given by

$$\cos\phi = \frac{-p^{*2}[1 + (k^2/m_A^2)\cos^2\theta^*] + k^2/4}{(\{p^{*2}[1 + (k^2/m_A^2)\cos^2\theta^*] + k^2/4\}^2 - k^2 p^{*2} \cos^2\theta^* (1 + k^2/m_A^2))^{1/2}}, \quad (4.6)$$

where θ^* is the angle between the τ and the parent Higgs boson A , measured in the rest frame of A . Since the Higgs boson decays isotropically, the ϕ distribution can easily be generated from Eq. (4.6) by choosing random values for $\cos\theta^*$. The result is shown in Fig. 5 for three different combinations of m_H and m_P . Obviously the distribution is always strongly peaked at its minimal possible value

$$\cos\phi_{\min} = \frac{-p^{*2} + \frac{1}{4}k^2}{p^{*2} + \frac{1}{4}k^2}; \quad (4.7)$$

$\phi_{\min}$ can thus be determined quite precisely already from a small number of events. Note that $\phi_{\min}$ is much more sensitive to the mass of the Higgs boson that decays into the τ leptons than to the mass of the second Higgs boson; the shapes of the curves for $(m_H, m_P) = (10, 20)$ GeV and $(10, 30)$ GeV differ much less than the curve for $(20, 20)$ GeV; a second, small peak originating from $P \rightarrow \tau^+\tau^-$ decays is probably unmeasurable in the first two cases. If $s \gg m_P^2 - m_H^2$ and $m_A^2 \gg 4m_{\tau}^2$ one finds indeed that $\cos\phi_{\min} \simeq 1 - 8m_A^2/s$ becomes independent of the "spectator" Higgs boson.

Unfortunately, the τ signal almost vanishes if $\omega < 1$. The reason is that in this region even a light Higgs boson decays into $c\bar{c}$ rather than $\tau^+\tau^-$; see Eq. (4.2). The branching ratio into τ leptons is sizable only if ω is close to one, where the total HP cross section is small. As a result, one has $\sigma(e^+e^- \rightarrow HP \rightarrow \tau^+\tau^- X) \lesssim 0.05$ pb if $\omega < 1$

at $\sqrt{s} = 60$ GeV. We have already briefly mentioned that existing models⁶ predict $\omega \gtrsim 1$; this is also intuitively reasonable, since $\omega < 1$ would imply that the ratio between top- and bottom-quark Yukawa coupling is even bigger than the ratio of their masses. Nevertheless, it would be nice to cover the area experimentally. Remember that only a small part of the region $\omega < 1$ is

FIG. 5. The differential cross section $d\sigma(e^+e^- \rightarrow HP \rightarrow \tau^+\tau^- X)/d\phi$ at $\sqrt{s} = 60$ GeV, where ϕ is the opening angle of the τ pair in the laboratory frame. The pairs of numbers in parentheses give (m_H, m_P) in GeV.

excluded by the $\Upsilon \rightarrow A\gamma$ bound; see Fig. 1.

One possible signal emerges when one of the two Higgs bosons decays into a $c\bar{c}$ pair, giving rise to two well-separated jets (the opening angle would be very similar to the one shown in Fig. 5). The second Higgs boson would then decay into $c\bar{c}$ or $b\bar{b}$ pairs. In the former case, one has four well-defined jets. The background is now of order $(\alpha_{\text{em}}\alpha_s)^2$ and, therefore, much bigger than in the case of the τ signal. Nevertheless, one should be able to extract the signal, since it produces peaks both in the opening angles and the invariant masses of the two jet pairs. Also, in 75% of the cases at least one c quark will decay semileptonically, which one might use to further reduce the background. Finally, since H and P decay independently, their two decay planes will be uncorrelated, whereas QCD four-jet events tend to be planar.³⁰

If, however, one or both Higgs bosons decay into $b\bar{b}$ pairs, the situation is somewhat different, since the b quarks might be rather slow. Furthermore, in many cases the Higgs-boson masses are not much bigger than $2m_b$; from Eq. (4.6) we see that in this case the opening angle between the b and the $\bar{b}$ will be small, so that only one broad jet might result from the $b\bar{b}$ system. In order to suppress QCD backgrounds one might then have to require that one b quark decay semileptonically, producing a sufficiently isolated charged lepton. Alternatively, one might try to reconstruct the invariant mass of the $b\bar{b}$ system; this will only be possible if both b quarks, as well as the c quarks produced from the b quarks, decay fully hadronically, the probability of which is only about 25%. Finally, in some regions of parameter space the event rate might be large enough to try direct b tagging, where one normally requires³¹ the b quark to hadronize into a B meson, which then has to decay into a charged D^* meson (branching ratio $\approx 20\%$). Which (if any) of these methods can be used to extract the $c\bar{c}b\bar{b}$ or $4b$ signal obviously depends crucially on detector properties and thus cannot be decided by us.

V. SUMMARY AND CONCLUSION

In this paper the associate production of a scalar and pseudoscalar Higgs boson at the e^+e^- collider TRISTAN was discussed. We worked within the framework of minimal supersymmetry, but most of our findings should also be valid in more general cases.

We first discussed the bounds that emerge within this model from the absence of monochromatic photons in Υ decays. We found that, since the model predicts $m_H \leq m_P$, $\Upsilon \rightarrow H + \gamma$ usually gives a stronger bound than $\Upsilon \rightarrow P + \gamma$, but is also more strongly affected by uncertainties related to large QCD corrections. We also briefly discussed uncertainties from bound-state effects and Coleman-Weinberg corrections to the scalar potential and concluded that the existing bound might not be absolutely reliable.

We then proceeded to show that TRISTAN can indeed cover a substantial region of parameter space. The

big advantage of the energy increase from PETRA to TRISTAN was stressed: The total cross section for HP production at $\sqrt{s} = 60$ GeV can be larger than 1 pb; most background cross sections should be smaller at the same time. While this cross section is probably too small to produce a measurable increase in R , the ratio of hadronic to QED-muonic total cross sections, it might be sufficiently large to produce a sizable number of characteristic events.

The most prominent signature is the $\tau^+\tau^-$ decay of one (or both) of the Higgs bosons. This is the dominant decay mode for a large part of the region where the total cross section exceeds 0.1 pb and where the $b\bar{b}$ mode is kinematically suppressed, provided that the ratio ω of vacuum expectation values is bigger than one as is usually expected. The event topology is a τ pair in one hemisphere and another τ pair or hadronic jets in the other. Since the τ 's are highly relativistic in most cases, one should be able to reconstruct the opening angle between them from their detected decay products; since the distribution of this opening angle has a strong Jacobian peak at its kinematic minimum, already a rather small number of events should allow the determination of the mass of the correspondings Higgs boson with an accuracy of a few GeV.

In the theoretically disfavored case $\omega < 1$, $c\bar{c}$ and/or $b\bar{b}$ decays of the Higgs bosons will have to be used. If both H and P decay into $c\bar{c}$ ($m_{H,P} \lesssim 2m_b$ or $\omega \ll 1$), the signature is four-jet events with double peaking in invariant mass and opening angle; in this case the QCD background should be manageable. If one or both Higgs bosons decay into $b\bar{b}$ pairs, phase-space constraints can easily lead to situations where one does no longer see four well-separated jets. In this case some properties of b decays will probably have to be used in order to reduce QCD backgrounds. Again, this problem should be less severe in the theoretically favored case $\omega > 1$, since here one would have to rely on the four- b -quark signature only if P and H are both quite heavy, which should increase the chance of getting more jetty events. However, definite conclusions about the four-quark signals clearly will have to wait until detailed Monte Carlo studies including detector efficiencies have been performed.

Nevertheless, it seems fair to say that with an integrated luminosity of 50–100 pb⁻¹ almost the whole kinematically accessible region can be probed. But even with a smaller integrated luminosity and based on the τ signature alone TRISTAN will already have a unique opportunity to shed some light on the mysteries surrounding electroweak symmetry breaking.

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