Production of pseudoscalar Higgs bosons in $e\gamma$ collisions

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We investigate the production of a pseudoscalar Higgs boson A^0 using the reaction $e\gamma\to eA^0$ at an $e\bar{e}$ collider with a center-of-mass energy of 500 GeV. Supersymmetric contributions are included which provide a substantial enhancement to the cross section for most values of the symmetry-breaking parameters. We find that, despite the penalty incurred in converting one of the beams into a source of backscattered photons, the $e\gamma$ process is a promising channel for the detection of the A^0.

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I. INTRODUCTION

The Higgs sector in supersymmetric extensions of the standard model contains charged Higgs bosons as well as additional neutral Higgs bosons [1]. Among the latter is a pseudoscalar particle usually denoted A^0. In this paper, we calculate the production cross section for the A^0 in the process $e\gamma\to eA^0$. Contributions to this process arise from triangle and box diagrams. The triangle contributions consist of diagrams in which the A^0 and photon are on-shell external particles and the remaining particle is a virtual photon or Z^0 in the t channel. Since $t=0$ is in the physical region, the photon pole contribution dominates the Z^0 pole contribution in this set of diagrams [2]. Moreover, because of the off-diagonal structure of the A^0 couplings to other bosons, the particles in the loop are either quarks, leptons, or charginos. Here, we present the top quark, bottom quark, τ lepton and the two chargino contributions to the photon pole amplitude.

The box diagrams have a more complex particle structure, with leptons, charginos, neutralinos, and scalar leptons in the loops. Like the Z^0 pole, these diagrams are nonsingular at $t=0$, and should not contribute a sizable correction to the photon pole terms. They are not included in the present calculation.

II. THE CROSS SECTION FOR A^0 PRODUCTION

The amplitude for the production of an A^0 of momentum k' and an e of momentum p' in the collision of an e of momentum p and a γ of momentum k and polarization $\varepsilon_\lambda(k)$ by the exchange of a γ in the t channel is

$$\mathscr{M}=\frac{4i\alpha^2}{\sin\theta_W m_W}\bar{u}(p')\gamma_\mu u(p)\frac{\mathscr{A}_\gamma(t)}{t}\varepsilon_{\mu\nu\alpha\beta}\varepsilon_\nu(k)\times(p-p')_\alpha k_\beta, \tag{1}$$

where $t=-(p-p')^2$, and

$$\begin{aligned}\mathscr{A}_\gamma(t)=[&-3(\tfrac{2}{3})^2 m_t^2\cot\beta C_0(t,m_A^2,m_t^2)\\&-3(-\tfrac{1}{3})^2 m_b^2\tan\beta C_0(t,m_A^2,m_b^2)\\&-(-1)^2 m_\tau^2\tan\beta C_0(t,m_A^2,m_\tau^2)\\&+2m_W m_1 g_{11}C_0(t,m_A^2,m_1^2)\\&+2m_W m_2 g_{22}C_0(t,m_A^2,m_2^2)].\end{aligned} \tag{2}$$

Here, m_t and m_b are the top and bottom quark masses, m_τ is the τ lepton mass, m_1 and m_2 are the chargino masses, m_W is the W mass, and $\tan\beta$ is a ratio of vacuum expectation values [1]. The chargino coupling constants g_{11} and g_{22} depend on the elements of two 2×2 unitary matrices U and V which diagonalize the chargino mass matrix X, where [3]

$$X=\begin{pmatrix}M & \sqrt{2}m_W\sin\beta\\ \sqrt{2}m_W\cos\beta & \mu\end{pmatrix}, \tag{3}$$

and which are chosen to ensure that m_1 and m_2 are positive. For illustrative purposes, we assume that the symmetry-breaking parameters M and μ are real and consider two cases: $M\mu>m_W^2\sin2\beta$ and $M\mu<m_W^2\sin2\beta$. The couplings in these cases are

$$g_{11}=\frac{m_W}{m_1^2-m_2^2}(m_2+m_1\sin2\beta),$$

$$g_{22}=-\frac{m_W}{m_1^2-m_2^2}(m_1+m_2\sin2\beta), \tag{4}$$

for $M\mu>m_W^2\sin2\beta$, and

$$g_{11}=\frac{m_W}{m_1^2-m_2^2}(-m_2+m_1\sin2\beta),$$

$$g_{22}=-\frac{m_W}{m_1^2-m_2^2}(-m_1+m_2\sin2\beta), \tag{5}$$

for $M\mu<m_W^2\sin2\beta$. Notice that these couplings are symmetric in m_1, m_2 and, unlike the A^0-top coupling, there is no

enhancement factor of $m_{1,2}/m_W$ [4]. We take $m_1 \geq m_2$. Because of the reality of M and μ, m_1 and m_2 in Eqs. (4) and (5) are subject to certain constraints discussed below [5]. The scalar function $C_0(t,m_A^2,m^2)$ is [6]

$$C_0(t,m_A^2,m^2)=\frac{1}{i\pi^2}\int d^4q \frac{1}{(q^2+m^2)[(q+p-p')^2+m^2][(q+p-p'+k)^2+m^2]}. \tag{6}$$

Since one of the external particles is a photon, this function can be expressed in terms of inverse trigonometric or hyperbolic functions [1,7] as

$$C_0(t,m_A^2,m^2)=\frac{1}{(t-m_A^2)}\left[C\left(\frac{m_A^2}{m^2}\right)-C\left(\frac{t}{m^2}\right)\right], \tag{7}$$

where

$$C(\beta)=\int_0^1 \frac{dx}{x}\ln[1-\beta x(1-x)-i\varepsilon] \tag{8}$$

$$=\begin{cases} 2\left[\operatorname{arcsinh}\left(\sqrt{-\frac{\beta}{4}}\right)\right]^2, & \beta\leq 0, \\ -2\left[\arcsin\left(\sqrt{\frac{\beta}{4}}\right)\right]^2, & 0\leq\beta\leq 4, \\ 2\left[\operatorname{arccosh}\left(\sqrt{\frac{\beta}{4}}\right)\right]^2-\frac{\pi^2}{2}-2i\pi\ \operatorname{arccosh}\left(\sqrt{\frac{\beta}{4}}\right), & \beta\geq 4. \end{cases} \tag{9}$$

The cross section is given by

$$\frac{d\sigma(e\gamma\to eA_0)}{d(-t)}=\frac{1}{64\pi s^2}\sum_{\text{spin}}|\mathscr{M}|^2, \tag{10}$$

and we have

$$\sum_{\text{spin}}|\mathscr{M}|^2=\frac{\alpha^4}{\sin^2\theta_W m_W^2}(s^2+u^2)\frac{|\mathscr{A}_\gamma(t)|^2}{(-t)}, \tag{11}$$

where $s=-(p+k)^2$ and $u=-(p'-k)^2$. The presence of the $1/t$ in Eq. (11) means it is necessary to introduce a cutoff in the calculation of the total cross section. One approach to obtain a finite cross section is to use the effective photon or Weizsäcker-Williams approximation for the exchanged photon [8]. Here, we integrate the exact amplitude and impose an angular cutoff. The expression for the total cross section is

$$\sigma_{e\gamma\to eA^0}(s)=\frac{\alpha^4}{64\pi\sin^2\theta_W m_W^2}\int_{\eta(s-m_A^2)}^{(s-m_A^2)}\frac{dy}{y}\left(2-2\frac{(m_A^2+y)}{s}+\frac{(m_A^2+y)^2}{s^2}\right)|\mathscr{A}_\gamma(-y)|^2, \tag{12}$$

where η is an angular cutoff. We investigated the effect of varying $\eta=\sin^2(\theta_{\min}/2)$ by comparing the standard model cross section with and without the Z^0 exchange. For $\theta_{\min}$ as large as $\pi/6$, the Z^0 contribution is only 3%–4% of the total. The result scales approximately as the logarithm of η, and we use $\eta=10^{-5}$ in the figures.

To complete the calculation of the cross section for the $e\gamma$ process, it is necessary to fold the cross section, Eq. (12), with the distribution $F_\gamma(x)$ of backscattered photons having momentum fraction x [9] to obtain

$$\sigma_T=\frac{1}{s}\int_{m_A^2}^{0.83s} d\hat{s}F_\gamma\left(\frac{\hat{s}}{s}\right)\sigma_{e\gamma\to eA^0}(\hat{s}), \tag{13}$$

with $\hat{s}=xs$. Here, we have taken the usual upper limit on the allowed x value, $x=0.83$.

This cross section is plotted in Fig. 1 for $M\mu>m_W^2\sin2\beta$ and in Fig. 2 for $M\mu<m_W^2\sin2\beta$. The dashed line in each panel is the contribution from the top and bottom quarks and the τ lepton. For large $\tan\beta$, the τ contribution is important. This is illustrated in the $\tan\beta=20$ panel of Fig. 1, where the dot-dashed line is the contribution from the top and bottom quarks. In Fig. 1, the solid lines are $m_1=250$ GeV and m_2 the largest value consistent with the constraint $(m_1-m_2)\geq m_W\sqrt{2(1+\sin2\beta)}$, which is needed to ensure that M and μ are real. Similarly, in Fig. 2, the solid lines correspond to $m_1=250$ GeV and m_2 the largest value consistent with $(m_1-m_2)\geq m_W\sqrt{2(1-\sin2\beta)}$. Unlike the $M\mu>m_W^2\sin2\beta$ case, when $M\mu<m_W^2\sin2\beta$, it is possible for m_1 and m_2 to be equal for $\tan\beta=1$ provided the $m_1,m_2\geq m_W$. These values of m_1 and m_2 are within the range of chargino masses found in studies of minimal supersymmetric models [10]. In most cases, the inclusion of the chargino contribution leads to a significant increase in the cross section, especially for the larger values of $\tan\beta$.

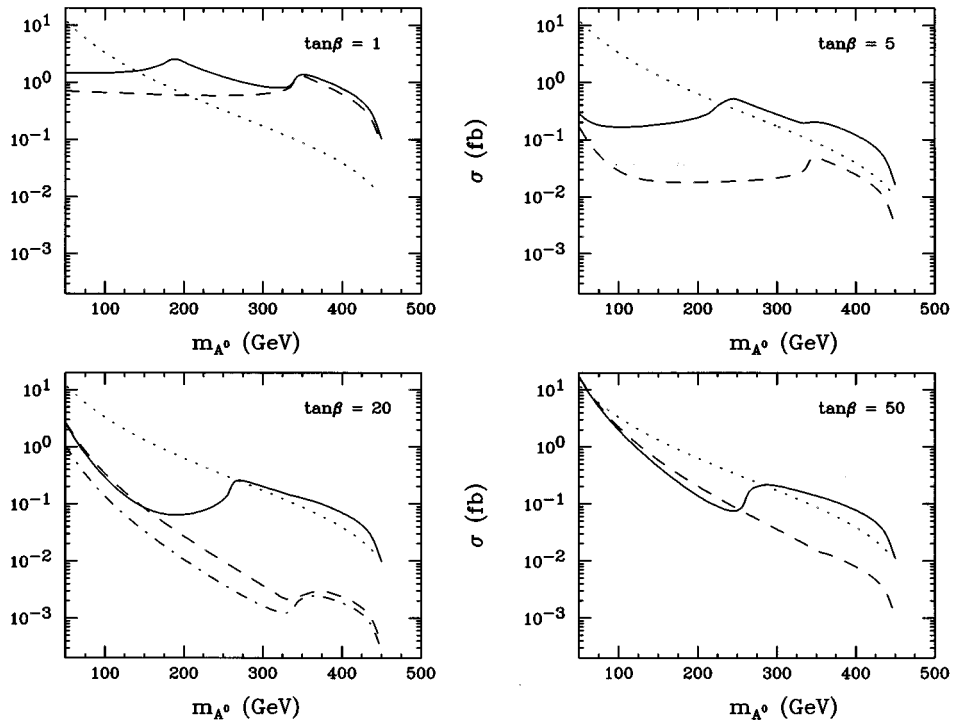

FIG. 1. Cross sections for the production of A^0 are shown for various values of $\tan\beta$ and an $e\bar{e}$ center-of-mass energy of 500 GeV when $M\mu > m_W^2 \sin 2\beta$. In each case, the solid line corresponds to chargino masses $m_1 = 250$ GeV and m_2 the largest value consistent with the restriction $(m_1 - m_2) \geq m_W\sqrt{2(1+\sin 2\beta)}$. The dashed line is the standard two Higgs doublet contribution without charginos, and the dot-dashed line in the $\tan\beta = 20$ panel is the two Higgs doublet result without the τ contribution. The dotted lines are the cross sections for the production of a background $b\bar{b}$ with invariant mass m_A. In these graphs, the angular cutoff η is taken to be 10^{-5}.

To assess the observability of this process, we assume that the dominant A_0 decay is $A_0 \to b\bar{b}$. For $m_A < 2m_t$, this ignores some contribution from chargino pair decay, but this is relatively small since in all but one of the examples we consider, the lowest chargino mass exceeds ~ 120 GeV. Even above the top threshold, $b\bar{b}$ decay dominates when $\tan\beta \sim 20$ [11]. The dotted lines in Figs. 1 and 2 are the cross sections for the direct production of a background $b\bar{b}$ of invariant mass m_A in $e\gamma$ collisions subject to an angular cut on the b and $\bar{b}$ direction relative to that of the incident photon in the $e\gamma$ center of mass. We find that an angular cut of $|\cos\theta| < 0.98$ on both the b and the $\bar{b}$ reduces the background $b\bar{b}$ cross section by about a factor of 10 while leaving the $b\bar{b}$ signal from A^0 decay essentially unchanged. More restrictive cuts on the b and $\bar{b}$ angles can further suppress the background, but at the expense of a significant decrease in the signal [12]. The cut shown appears to be optimal.

III. DISCUSSION

We would like to point out that the $e\gamma$ cross sections calculated here are very likely to be much larger than those

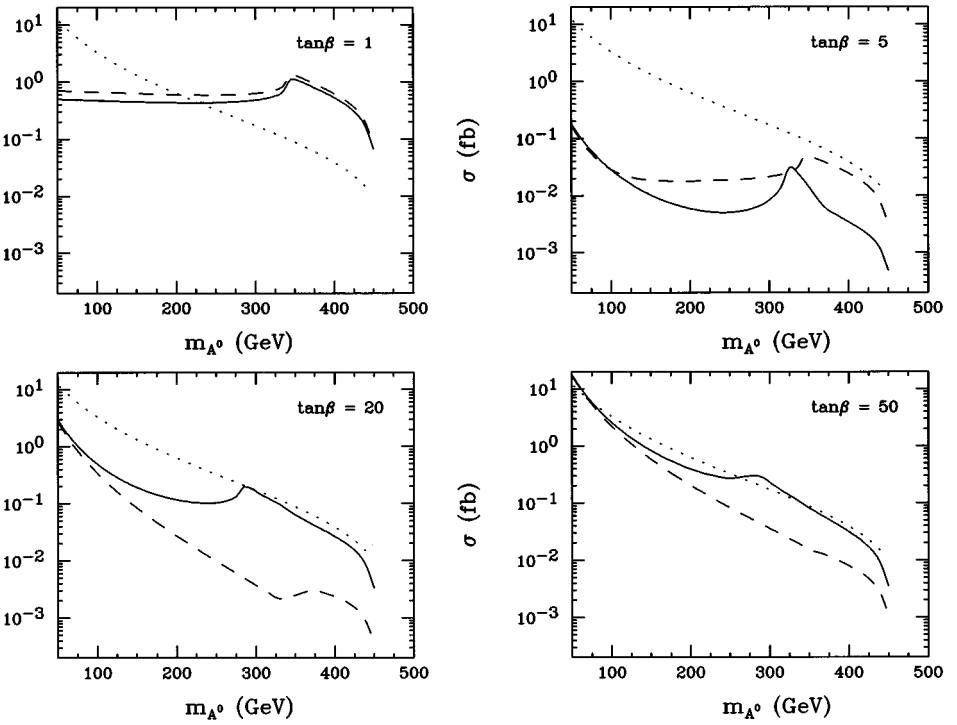

FIG. 2. Same as Fig. 1 for $M\mu < m_W^2 \sin 2\beta$. In this case, m_1 and m_2 satisfy the condition $(m_1 - m_2) \geq m_W\sqrt{2(1-\sin 2\beta)}$.

of the related process $e\bar{e}\to\gamma A^0$ at 500 GeV. We have checked this for the production of the standard model Higgs boson using the complete (standard model) calculation of $e\bar{e}\to\gamma H^0$ [7] and the photon pole contribution to $e\gamma\to eH^0$. At an $e\bar{e}$ center-of-mass energy of 500 GeV, we find the cross section $\sigma(e\bar{e}\to\gamma H^0)$ for the production of a 200 GeV H^0 is 0.08 fb, whereas $\sigma(e\gamma\to eH^0)=5.9$ fb for the same Higgs-boson mass.

This enhancement is implicit in a previous calculation of scalar Higgs-boson production [8] where the Weizsäcker-Williams approximation is used for the t-channel photon together with the on-shell $H\to\gamma\gamma$ amplitude. This is essentially equivalent to setting $y=0$ in the parentheses of Eq. (12) and using m_e^2 as the cutoff in the remaining integral [13]. Our comparison of the approximate results of Ref. [8] with an exact calculation suggests that the Weizsäcker-Williams approach tends to overestimate the cross section. Apart from minor variations depending on how the calculation is performed, it is nevertheless true that the t-channel cross section is substantially larger than its s-channel counterpart.

The production of A^0 has also been investigated in $\gamma\gamma$ collisions [14]. In that case, the background arises from the process $\gamma\gamma\to b\bar{b}$ and it is effectively suppressed by imposing an angular cut. With our choice of chargino masses, a comparison of the $\tan\beta=20$ cross sections in $e\gamma$ production and $\gamma\gamma$ production [14] reveals a larger signal in the $e\gamma$ mode. Both channels are likely to be important in searches for the A^0.

To the extent that the photon pole contribution can be isolated, this method of searching for the A^0 has the advantage that the contributions from supersymmetry are significant and limited to one type of supersymmetric particle. Should one observe a cross section larger than any standard model prediction, the case for the presence of chargino contributions is rather strong.

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