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Light Neutral Higgs Bosons at the Low Energy $\gamma\gamma$ Collider

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ABSTRACT

A light neutral Higgs boson in the framework of the general two Higgs doublet model (2HDM) is not excluded by existing data. We point out that it can be looked for at the proposed low energy $\gamma\gamma$ collider. Failure to detect one may lead to important limits on the parameters of the general 2HDM.

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While the Standard Model (SM) Higgs scalar as well as the MSSM neutral Higgs particles have been constrained by LEP1 data to be heavier than 65.2 GeV, and 40–50 GeV, respectively, the general two Higgs doublet model (2HDM) may yet accomodate a very light ($\lesssim 40$ GeV) neutral scalar h or a pseudoscalar A as long as $M_h + M_A \gtrsim M_Z$ [1, 2]. The interesting case of very light ($\sim$ few GeV) Higgs particles has been studied in dedicated experiments, for example, in the Wilczek process [3]. Unfortunately, limits are not decisive, especially due to large theoretical uncertainties in QCD and in relativistic corrections. There is some hope though that better limits may be obtained by exploring the Yukawa process ($Z \to f\bar{f}h/A$) in the existing LEP1 data [4] or in improved $(g-2)_\mu$ measurements [5].

The $\gamma\gamma$ option at the Next Linear Collider, on the other hand, may provide an excelent opportunity to search for a *very light neutral Higgs particle*. We focus here on the resonant production of a very light neutral Higgs particle at the low energy $\gamma\gamma$ collider, suggested as a test machine for the NLC [6].

The general 2HDM is charactarized by five (Higgs) masses and two parameters (angles): α and β [7]. We consider here the phenomenologically appealing version, where the neutral components of the two doublets $\phi_{1,2}$ (with vacuum expectation values $v_{1,2}$) couple exclusively to the $I_3 = \pm 1/2$ fermion fields. Tree level flavour changing neutral currents, then, vanish identically. Such an assumption could lead naturally to a large value for the ratio $\tan\beta \equiv v_2/v_1$ ($\sim m_t/m_b \gg 1$) and, thus, to an enhanced coupling of the light scalar (pseudoscalar) to the *down*-type quarks and the charged leptons, while suppressing the coupling to the *up*-type quarks.

In addition to the above, the extension to the 2HDM results in significant modification in the scalar–vector boson sector of the theory. The canonical Higgs boson production mechanism, namely the Bjorken process $Z \to Z^* h$, now proceeds with a rate proportional to $\sin^2(\alpha - \beta)$. Negative results at LEP1 thus imply that $\sin^2(\alpha - \beta) < 0.1$ if $M_h \lesssim 50$ GeV. More than this, a strikingly new feature is that the Z can now couple to a pair of nonidentical spin-0 objects, leading to new Higgs particle production mechanisms. Of particular interest[1] is the process $Z^* \to h + A$, with a rate $\propto \cos^2(\alpha - \beta)$. A lack of such events at LEP1 can then be translated to a constraint in the three-dimensional $(M_h, M_A, \alpha - \beta)$ parameter space [2]. For a given (M_h, M_A) combine, this obviously translates to an relatively strong upper limit on $\cos^2(\alpha - \beta)$. The two constraints are thus complementary to each other. In addition, one has also to consider the fact that a non-trivial Higgs sector may lead to additional contribution to the Z width, even if the new decay channels cannot be identified over the SM background. Yet, a light Higgs pair ($M_h + M_A \lesssim 70$ GeV) may still be accomodated [2].

In this Letter, we concentrate on the scenario wherein either h or A is very light [8]. While, for a light h, this clearly warrants that $\alpha \simeq \beta$, it is not necessary if only A is light and $M_h > m_Z - M_A$. However, in order to reduce the number of parameters and thus simplify the analysis, we shall not only impose this constraint, but rather promote it to an exact equality. Since we propose to use charged lepton decay modes as our signal, we further restrict ourselves to the scenario with a large $\tan\beta$ ($\gtrsim 20$). Note that, for $M_{h,A} < 10$ GeV, this parameter is not

[1]Note that there are no tree level ZZA or W^+W^-A vertices in this theory.

yet constrained by LEP1 data. Although, for $M_{h,A} \gtrsim 10$ GeV, non-observation of the Yukawa process ($Z \to b\bar{b}h/A$) constrains $\tan\beta$ to be below 10(5) for scalar (pseudoscalar) [9], the same process is unlikely to be as efficient for lower Higgs boson masses. Such an analysis is currently in progress though [4]. Low energy data like those on the muon anomalous magnetic moment still allow $\tan\beta \sim 20$ or higher for $M_h \gtrsim$ 2–3 GeV [10, 12] (see later discussion).

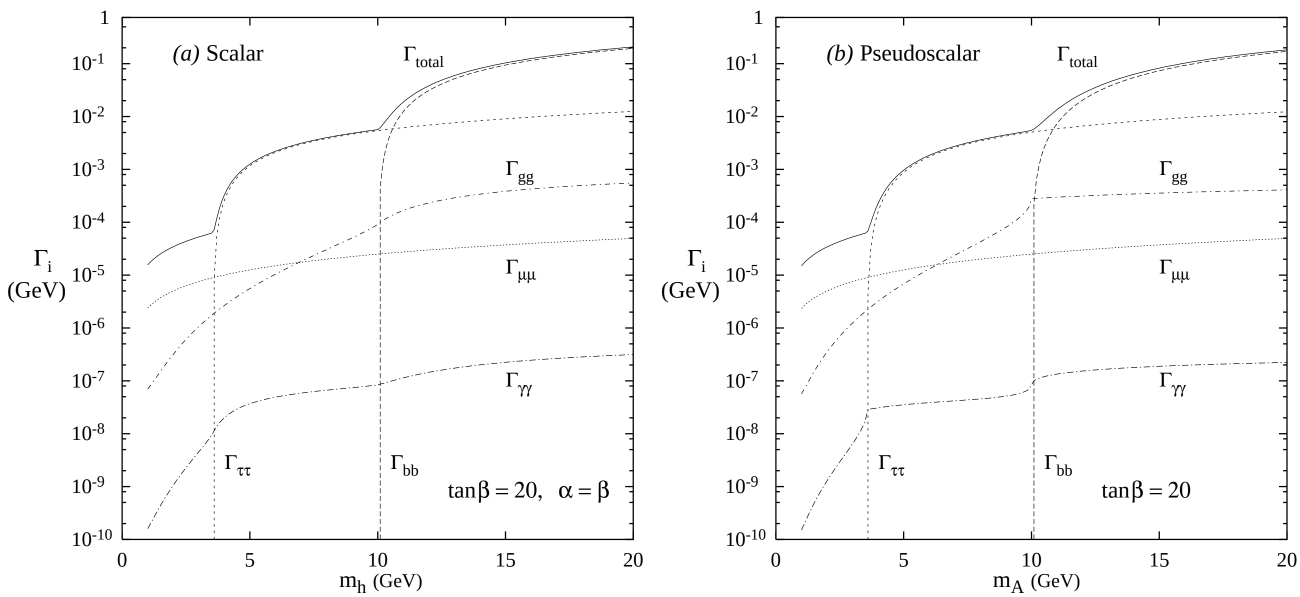

Figure 1: *The partial and total decay widths for* $\tan\beta = 20$*:* (a) *scalar* h *(*$\alpha = \beta$*), and* (b) *pseudoscalar* A*.*

In Fig.1, we present relevant for our analysis partial widths of h and A (for $\tan\beta = 20$) obtained under the above hypothesis. To the leading order, all the widths shown in the figure scale as $\tan^2\beta$. Note that the fermionic branching fractions for the two cases follow a very similar pattern, and the only noticebale difference occurs in the 2-photon and the 2-gluon decay modes.

Resonant neutral (pseudo)scalar production may occur at a $\gamma\gamma$ collider through two photon fusion at one loop. While only the charged fermion loops contribute for A, the $\gamma\gamma h$ vertex would, in general, receive corrections from $W^\pm$ and $H^\pm$ loops as well. However, for $\alpha = \beta$ the $W^\pm$ contribution vanishes identically. In the same limit, the H^+H^-h vertex is proportional to $(gm_Z/4\cos\theta_W)\ \sin(4\beta)$, where g is the weak coupling constant. Clearly, this vertex becomes progressively weaker as $\tan\beta$ increases beyond 20. Moreover, this contribution weakens further as M_{H^+} increases. Since we assume that the charged scalars are indeed heavy, this contribution can be safely neglected for the purpose of our study.

The cross section for the basic process $\gamma\gamma \to h \to f\bar{f}$ (where we specify f to be τ or μ as the most important decay modes) is given by[2]

$$\sigma_{\gamma\gamma} = \frac{8\pi\Gamma(h \to \gamma\gamma)\Gamma(h \to f\bar{f})}{(s_{\gamma\gamma} - M_h^2)^2 + \Gamma_h^2 M_h^2}\,(1 + \lambda_1\lambda_2)\ , \tag{1}$$

[2] A similar expression holds for the pseudoscalar.

where λ_i are the mean helicities of the photon beams and the rest of the symbols carry their usual meaning. In order to calculate the total cross section $\sigma^h_{ee}(f\bar{f})$, we need to fold the above with the appropriate photon spectrum. Thus, for an e^+e^- center of mass of $\sqrt{s_{ee}}$, the differential cross section is given by

$$\frac{\mathrm{d}\sigma^h_{ee}}{\mathrm{d}s_{\gamma\gamma}} = \int_{x_{\min}}^{x_{\max}} \frac{\mathrm{d}x_1}{x_1 s_{ee}} \, \sigma_{\gamma\gamma} \; f(x_1) \; f\left(\frac{s_{\gamma\gamma}}{s_{ee}x_1}\right), \tag{2}$$

where x_1 (and $x_2 \equiv s_{\gamma\gamma}/s_{ee}x_1$) are the momentum fractions of the initial electrons carried by the photons, and $x_{\min} = s_{\gamma\gamma}/s_{ee}/x_{\max}$. The photon spectrum $f(x)$, resulting from Compton backscattering electrons on an intense laser light, depends [11] on the helicity of the initial electrons λ_e, initial laser beam circular polarization P_c and a machine parameter z that determines the maximum momentum carried by the photon ($x_{\max} = z/(1+z)$). Defining, for convenience, variables $r \equiv x/(1-x)$ and $y \equiv 1 - x + 1/(1-x)$, we have, for the photon spectrum and the mean helicity $\lambda_\gamma(x)$,

$$\begin{aligned} \frac{\mathrm{d}n(x)}{\mathcal{N}\mathrm{d}x} &= y + 4\frac{r}{z}\left(\frac{r}{z} - 1\right) - \lambda_e \, P_c \, r(2-x)\left(\frac{2r}{z} - 1\right) \;, \\ \lambda_\gamma(x) &= \left(\frac{\mathrm{d}n(x)}{\mathcal{N}\mathrm{d}x}\right)^{-1} \left[-\lambda_e r \left\{1 + (1-x)\left(\frac{2r}{z} - 1\right)^2\right\} + P_c \, y \left(\frac{2r}{z} - 1\right)\right] \;, \end{aligned} \tag{3}$$

where $\mathcal{N}$ gives the normalization.

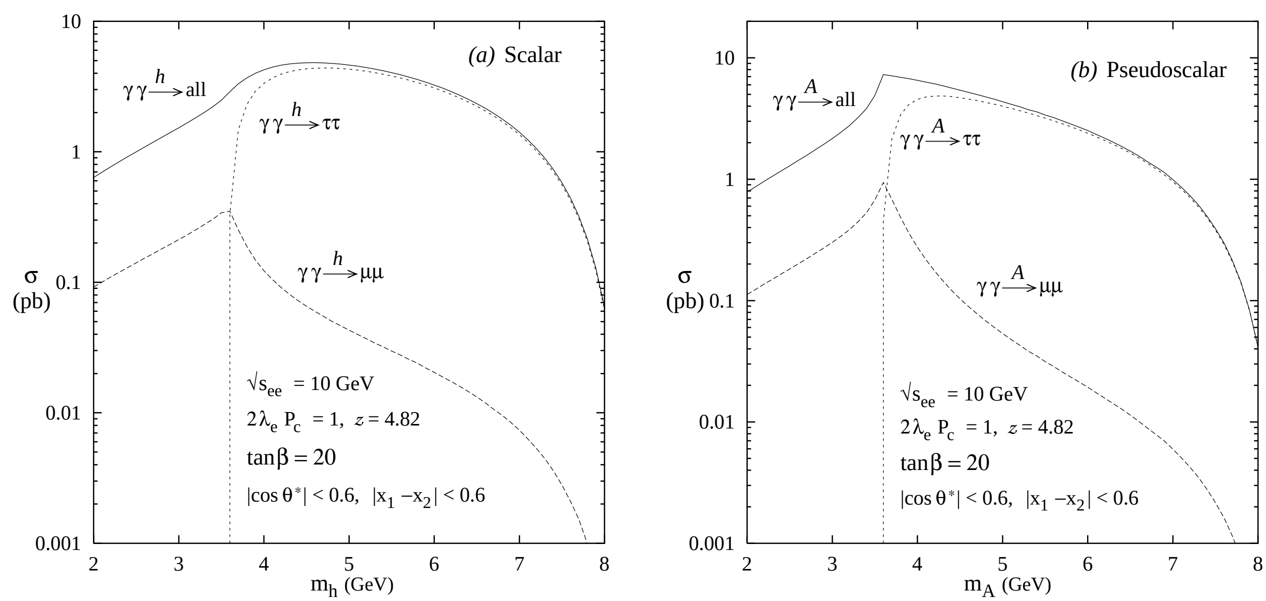

Figure 2: *The effective cross sections for the Higgs boson mediated process leading to $\mu^+\mu^-$ and $\tau^+\tau^-$ final states ($\tan\beta = 20$). Also shown is the cross section summed over all decay channels.* (a) *scalar h ($\alpha = \beta$) and* (b) *pseudoscalar A.*

We consider the resonant production of very light Higgs scalar (pseudoscalar) at e^+e^- NLC collider with energy $\sqrt{s_{ee}}$=10 GeV [6]. To maximize the photon energy, and yet avoid multiple rescattering or pair–creation [11], we choose $z = 2(\sqrt{2}+1) = 4.82$, and thus $\sqrt{s_{\gamma\gamma}^{\max}} \simeq 0.83\sqrt{s_{ee}} = 8.3$ GeV. Following ref. [11], we assume the 'broad' spectrum of photons with $2\lambda_e P_c = +1$. This has the advantage of being rather flat over $\sqrt{s_{\gamma\gamma}}$ and, more importantly, of favoring the $J_Z = 0$ state, the polarization state of Higgs scalar. To a very good approximation, $\sigma^h_{ee}(f\bar{f}) \propto \tan^2\beta$. In Fig.2, we present the cross sections for $\tan\beta = 20$. The cuts on the center of mass angle θ^* and the difference in the photon momentum fraction are motivated below.

The (orders of magnitude larger) background is mainly due to direct $f\bar{f}$ pair production, since at such low energy the resolved photon contributions are negligible. In Fig.3 we show the invariant mass distribution for the QED process. While the predominant contribution ($J_Z = 2$)

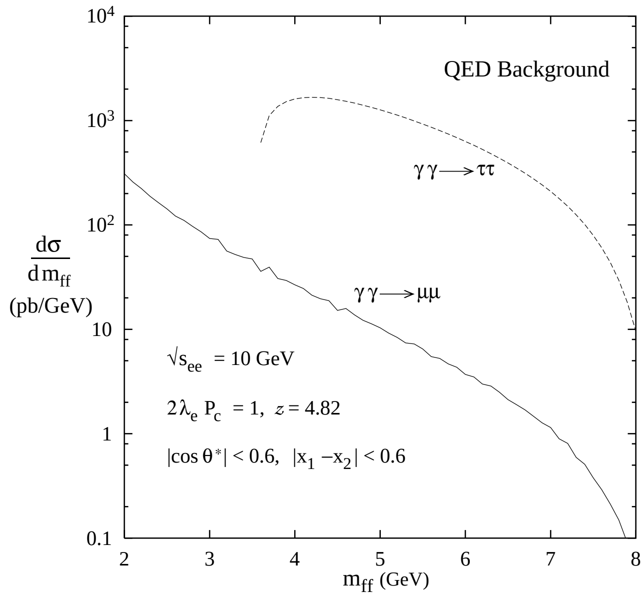

Figure 3: *The invariant mass distribution for the QED process $\gamma\gamma \to f\bar{f}$.*

is already reduced significantly by the choice for the spectrum, the forward/backward peaked $J = 0$ contribution is reduced by imposing a cut ($-0.6 < \cos\theta^* < 0.6$) on the CM scattering angle. As the signal is independent of θ^*, this cut eliminates only 40% of the Higgs events. It should be noted that this cut is more effective for the muonic channel than for the tauonic channel on account of the smaller mass of the muon. Although similar reduction can be obtained for the $\gamma\gamma \to \tau^+\tau^-$ by further restricting θ^*, this does not lead to an appreciable improvement in the signal to background ratio. Further improvements in this ratio can be made by restricting the total boost of the $f\bar{f}$ system, or equivalently, by restricting the difference in the momentum fractions carried by the two colliding photons. We find that $|x_1 - x_2| < 0.6$ is an optimal choice.

It is clear from eqn.(1), that the signal would have a a sharp peak in the $f\bar{f}$ invariant mass. The task then is to look for it over the continuous, but much larger, QED background. Since Γ_h (Γ_A) is tiny (see Fig. 1), in the event of infinite resolution in the invariant mass, the signal would be striking indeed! We adopt, though, a more realistic approach and smear the

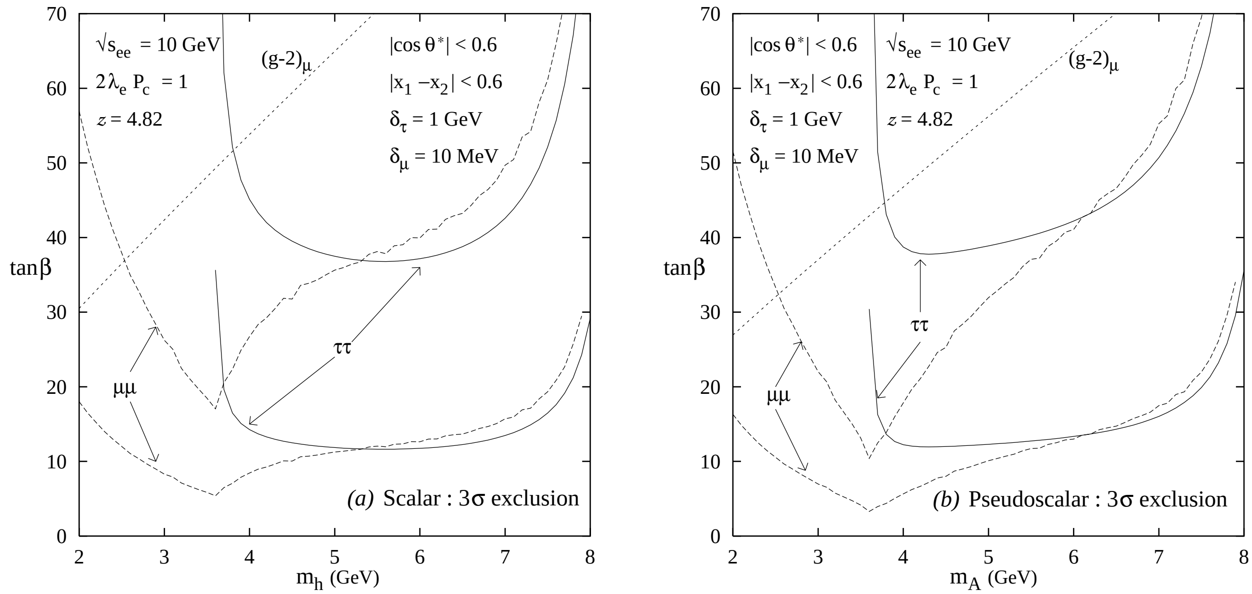

Figure 4: *The exclusion plots in the* (a) M_h–$\tan\beta$ *and* (b) M_A–$\tan\beta$ *planes that may be achieved at NLC$\gamma\gamma$ from either of $\mu\mu$– and $\tau\tau$–channels. Parameter space above the curves can be ruled out at 3σ (99.7% C.L.). The upper and lower sets are for integrated luminosity of 100 pb^{-1} and 10 fb^{-1} respectively. Also shown are the limits from the current data on $(g-2)_\mu$ under the assumption that the only non-SM contributions accrue from a light h/A.*

signal (as well as the background) profile with a gaussian resolution function [13]. Thus,

$$\frac{d\sigma^h_{ee}}{dm_{f\bar{f}}} = \frac{1}{\delta_f\sqrt{2\pi}} \int_{4m_f^2}^{x^2_{\rm max} s_{ee}} ds_{\gamma\gamma}\, \frac{d\sigma}{ds_{\gamma\gamma}}\; \exp\left(-\frac{(m_{f\bar{f}} - \sqrt{\sigma_{\gamma\gamma}})^2}{2\delta_f^2}\right). \tag{4}$$

For the experimental resolutions, we choose $\delta m_{f\bar{f}} = 2\delta_f$, with δ_μ=0.01 GeV and δ_τ=2 GeV. The 3σ (99.7% C.L.) exclusion plots in the $\tan\beta$–$M_{h/A}$ plane that may then be obtained using the $\mu^+\mu^-$ and the $\tau^+\tau^-$ final states are displayed in Fig.4. To be specific, we have adopted two representaive values for the integrated luminosity : 100 pb^{-1} and 10 fb^{-1}. It is interesting that though the $\tau^+\tau^-$ cross sections are typically larger, yet better limits are obtained from the muonic channel. The reasons are twofold : *(i)* as we have noted earlier, the angular cuts are more effective in eliminating the QED muons than the tau's and *(ii)* the invariant mass resolution for a $\tau^+\tau^-$ pair is expected to be much worse than that for the $\mu^+\mu^-$ pair. In the ideal case where $\delta_\tau \simeq \delta_\mu$, the bounds from the two channels would be similar. Note that above results are not expected to have large theoretical uncertainties in contrast to *e.g.*, the process [3, 12] $\Upsilon \to \gamma h(A)$ which, in principle, is sensitive to the same mass range.

Also displayed in Fig.4 are the bounds that can be obtained from the measurement [14] of the anomalous magnetic moment of the muon under the assumption that a light Higgs particle constitutes the only new source of contribution. Although the SM central value for this quantity depends on the evaluation of hadronic vacuum polarization, the difference is miniscule. For the curves above, we use Case A of ref. [15] as this provides the stricter bounds. A look at the figures tell us then that, even for the low luminosity version of NLC$\gamma\gamma$, the bounds obtainable from the experiment suggested here would be much stronger than the current ones. It should be borne in mind though that a substantial improvement in the experimental measurement of $(g-2)_\mu$ is in the offing [16] and if, in addition, theoretical errors could be reduced, this measurement could provide much stronger constraints.

To summarize, we consider the physics potential of high luminosity, low energy NLC$\gamma\gamma$ in the search for a possible light Higgs (pseudo)scalar in the general 2HDM. The limits obtainable with an integrated luminosity of 100 pb^{-1} would already be better for the mass range 3 GeV $\lesssim \mathrm{M_h} \lesssim$ 8 GeV than the existing limits from $(g-2)_\mu$. A luminosity of the order 10 fb^{-1}/y would of course lead to much more stringent bounds. The latter compare very well with those expected from a 20-fold improvement in $(g-2)_\mu$ measurement and do not depend upon any additional assumptions regarding the spectrum of new physics. Photon polarization is crucial though.

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