PHYSICAL REVIEW D VOLUME 38, NUMBER 11 1 DECEMBER 1988

Can the Higgs bosons of the minimal supersymmetric model be detected at a hadron collider via two-photon decays?

J. F. Gunion
Department of Physics, The University of California, Davis, California 95616
and Institute for Theoretical Physics, The University of California at Santa Barbara, Santa Barbara, California 93106

G. Gamberini and S. F. Novaes
Lawrence Berkeley Laboratory, The University of California, Berkeley, California 94720
(Received 21 June 1988)

Detection at a hadron collider of the neutral Higgs bosons of the minimal supersymmetric extension of the standard model is not straightforward. Their decays to WW and ZZ pairs are either phase-space forbidden or are suppressed or absent at the coupling-constant level; thus, detection of these neutral Higgs bosons in such modes will not be possible. Their detection in heavy-quark decay modes is also almost certainly impossible due to large QCD backgrounds. One possible detection mode (that proves useful for a standard-model Higgs boson with mass below $2m_W$ and $2m_{\mathrm{top}}$) is via decays to two photons. We investigate the branching ratio for such decays as a function of Higgs-boson mass for a selection of neutralino and chargino mass-matrix parameters. Except for a few limited regions of parameter space, we conclude that this mode of detection is probably not viable given the known two-photon continuum backgrounds.

I. INTRODUCTION

Probing the source of electroweak symmetry breaking is one of the most important goals of the next generation of colliders. In particular, we must ascertain the extent to which hadron colliders, such as the proposed Superconducting Super Collider (SSC) and Large Hadron Collider (LHC), can be used to carry out the necessary studies. While there is no general agreement as to the mechanism responsible for electroweak symmetry breaking, the standard-model (SM) scenario in which the observable experimental remnant is a single neutral Higgs boson, denoted by ϕ^0, suffers from problems of hierarchy and naturalness. Supersymmetry was introduced as a means for avoiding these problems.[1] The Higgs-boson sector of the minimal supersymmetric model (MSSM) has been studied in detail in Refs. 2–4. At least two doublets of Higgs fields are required in order to give masses to all the quarks and leptons and to guarantee the absence of anomalies. Five physical Higgs bosons emerge: a charged pair $H^\pm$, a light neutral scalar h^0, a heavy neutral scalar H^0, and a pseudoscalar A^0. A convenient summary of the properties of these Higgs bosons appears in Ref. 5. For the moment we merely remind the reader that all couplings and masses of the Higgs sector are determined by two parameters which we take to be

$$\tan\beta = v_2/v_1 \quad \text{and} \quad m_{H^\pm} , \tag{1.1}$$

where v_2 and v_1 are the vacuum expectation values of the Higgs fields which give mass to the up and down quarks, respectively, and $m_{H^\pm}$ is the mass of the charged Higgs boson. The pseudoscalar- and scalar-Higgs-boson masses are given in terms of these two parameters by

$$m_{A^0}^2 = m_{H^\pm}^2 - m_W^2 \tag{1.2}$$

and

$$m_{H^0,h^0}^2 = \tfrac{1}{2}\{ m_{A^0}^2 + m_Z^2 \pm [(m_{A^0}^2 + m_Z^2)^2 - 4m_Z^2 m_{A^0}^2 \cos^2 2\beta]^{1/2} \} . \tag{1.3}$$

These relations imply (i) $m_{H^0} \geq m_Z$, (ii) $m_{h^0} \leq m_{A^0}$, (iii) $m_{h^0} \leq m_Z |\cos 2\beta| \leq m_Z$, and (iv) $m_{H^\pm} \geq m_W$. Typically, we expect[2] $\tan\beta \gtrsim 1$; the nearer $\tan\beta$ is to 1 the lighter the h^0.

We also recall that there is no coupling of A^0 to WW or ZZ and that the coupling of H^0 to these VV channels is generally severely suppressed. Thus, the only Higgs boson in the MSSM with SM-like couplings to VV is the h^0. These facts have implications both for the decay and the production of these three neutral Higgs bosons. Consider first their decays. Since the h^0 is too light to decay to VV channels, while H^0 and A^0 have suppressed or zero coupling to VV, we find that all the Higgs bosons of the MSSM will decay predominantly to heavy-quark channels, unless non-SM channels become important. The above coupling pattern also implies that at a hadron collider VV fusion is never important for the production of the neutral Higgs bosons. First, the h^0, despite having large VV couplings, is constrained to be so light that gluon-gluon fusion via heavy-quark loops is dominant in the m_{h^0} range allowed by the theory. Second, the H^0 and A^0, which can be heavy enough that VV fusion could in principle have been an important production process, have suppressed or absent VV couplings and VV fusion processes can be ignored. Again, the dominant production process is gluon-gluon fusion via heavy-quark loops. The gg fusion cross sections for h^0, H^0, and A^0, as a function of the appropriate mass (m_{h^0}, m_{H^0}, or m_{A^0}, respectively), are all similar to that computed for a SM Higgs boson of the same mass.[3] In particular, the cross sections are very dependent upon the size of the top-quark mass. For $m_t = 40$ GeV and SSC energy of $\sqrt{s} = 40$ TeV, the SM Higgs-boson cross section typically passes below 1 pb at a mass above ~ 500 GeV, while if $m_t = 200$ GeV the cross section remains above 1 pb all the way out to a mass of order 1 TeV (Ref. 6).

In this paper, we wish to focus on detection of the neutral MSSM Higgs bosons at a hadron collider such as the SSC. We presume that their detection in heavy-quark decay channels is impossible due to the very large QCD backgrounds. A mode that has been demonstrated[7] to be useful in the case of the SM Higgs boson ϕ^0 is $\phi^0 \to \gamma\gamma$. As long as m_{ϕ^0} is above about 80 GeV and below $\min\{2m_W, 2m_t\}$, the ϕ^0 can be detected in its two-photon decay mode above the two-photon continuum background from $q\bar{q} \to \gamma\gamma$ and $gg \to \gamma\gamma$ (where the latter process occurs at one loop). In this region of m_{ϕ^0} the branching ratio for $\phi^0 \to \gamma\gamma$ is of order 10^{-3} or a little smaller. Detection requires excellent mass resolution in the $\gamma\gamma$ channel (of order 2%), but this is expected to be possible.[8] It is also necessary to be able to discriminate between jets and isolated photons at a level of 1 part in 10^4. In this regard, the techniques employed by the CDF Collaboration in their recent measurements of the high-p_T spectrum of isolated photons seem to have already achieved this level.[9] In considering use of this same mode for the MSSM neutral Higgs bosons, we must recompute the gg fusion cross section for the particular Higgs boson in question, including squark loops as well as quark loops, and incorporating all mixing-angle effects. Programs for this purpose were developed and some results given in Ref. 3. Of course, the $\gamma\gamma$ continuum backgrounds are unchanged. By combining these two ingredients we can determine what branching ratio is required for h^0, H^0, or A^0 detection via decays to $\gamma\gamma$.

In more detail, let us consider SSC energy and luminosity of $\sqrt{s} = 40$ TeV and $L_{\text{year}} = 10^4$ pb^{-1}. We compute the background from $q\bar{q} \to \gamma\gamma$ and $gg \to \gamma\gamma$ [the latter is included approximately by multiplying the former cross section by a factor of 2 (Ref. 10)] as a function of $M_{\gamma\gamma}$ within a 2% mass bin around $M_{\gamma\gamma}$:

$$\Sigma_{\gamma\gamma}(M_{\gamma\gamma}) = \int_{\Delta M_{\gamma\gamma} = 0.02 M_{\gamma\gamma}} \frac{d\sigma}{dM_{\gamma\gamma}} dM_{\gamma\gamma} \ . \tag{1.4}$$

In computing Σ we impose a cut of $\cos\theta^* \le 0.5$ upon the outgoing photons in their center of mass, where θ^* is defined with respect to the beam axis. This cut considerably reduces the $\gamma\gamma$ continuum background. We next compute the cross section for gluon-gluon fusion production of the Higgs boson in question (h^0, H^0, or A^0), multiplied by a factor of $\frac{1}{2}$ to account for the $\cos\theta^* \le 0.5$ decay-angle restriction. Of course, the top-quark mass and the masses of the squarks must be specified in order to carry out the gluon-gluon fusion calculation. Values for these masses are chosen as specified in the sections that follow. The resulting cross section as a function of Higgs-boson mass m_h is denoted by $\Sigma_h(m_h)$, where h can be h^0, H^0, or A^0. In order to achieve at least a nominal 5σ effect in the signal compared to the background we must have a two-photon branching ratio for the Higgs boson in question of more than

$$B_{\min}^{h \to \gamma\gamma}(m_h) = 5 \frac{[L_{\text{year}} \Sigma_{\gamma\gamma}(M_{\gamma\gamma} = m_h)]^{1/2}}{L_{\text{year}} \Sigma_h(m_h)} \ . \tag{1.5}$$

This function will be plotted for comparison with the actual $\gamma\gamma$ branching ratios of the three neutral Higgs bosons in the figures discussed below.

II. TWO-PHOTON DECAY WIDTHS

In Ref. 5 the formalism for computing the widths for the decay of the neutral MSSM Higgs bosons to two photons was presented. Related results appear in Ref. 11. In Appendix A we summarize the necessary formulas in a particularly convenient form. In Ref. 5 we used the two-photon decay widths of the MSSM Higgs bosons to compute their production cross sections in two-photon collisions at an e^+e^- collider. Here we shall present results for the widths themselves, and for the resulting branching ratios $B(h^0, H^0, A^0 \to \gamma\gamma)$, after including other channels that contribute to the decays. It is useful to review some of the salient features of the calculations. First, we remind the reader that one of the largest contributing loops to the two-photon decays of the SM Higgs boson is that containing the W boson. In the MSSM, this loop is only important in the case of h^0, which is the only Higgs boson with strong couplings to WW. In the absence of non-SM loop contributions the branching ratios of H^0 and A^0 to two photons would be too small to be utilized. However, there are new loop contributions. The most important are those from the chargino partners of the charged Higgs boson and the W, and the squark partners of the SM quarks. In addition, there are lepton loops and charged-Higgs-boson loops.

In order to compute the above non-SM particle loop contributions, we must specify several additional parameters of the MSSM. The chargino sector of the theory is determined by the gaugino mass M associated with the SU(2) subgroup, and a supersymmetric Higgs-boson parameter μ. (See Ref. 4 for a detailed discussion.) In most approaches to minimal supersymmetry it is possible to relate the SU(2) gaugino mass to the mass of the SU(3) gluino, $M_{\tilde{g}}$, by requiring that all gaugino masses are equal at some grand unification scale. Then, at the electroweak scale, all gaugino mass parameters can be expressed in terms of one of them. In particular,

$$M_{\tilde{g}} = \left[\frac{g_s^2}{g^2} \right] M \ , \tag{2.1}$$

where g_s is the strong coupling constant. In this work, we shall present results for three values of M; 50 GeV, 200 GeV, and 500 GeV, corresponding to gluino masses of roughly 200 GeV, 800 GeV, and 2 TeV, respectively. In addition, at each M value we shall consider a related selection of μ values,

$$\mu = -2M, -M/2, 0, M/2, 2M \ , \tag{2.2}$$

as a representative of the possibilities. Note, however, that at $M=500$, $\mu=0$ and at $M=50$, $\mu=100$, the lighter chargino mass is below 15 GeV for both values of $\tan\beta$ we shall consider ($\tan\beta=1.5$ and $\tan\beta=4$). This is also the case at $M=200$, $\mu=0$ when $\tan\beta=4$. Such a light chargino mass is in disagreement with experimental limits;[12] the associated curves are included in the figures that follow only to illustrate the effect of having an extremely light chargino. To fix the squark sector, we need only specify the squark masses. For simplicity, we have taken them all to be degenerate with mass of $M_{\tilde{q}}=500$ GeV. Sensitivity to this mass scale is small as long as the Higgs boson of interest has mass $\lesssim 2M_{\tilde{q}}$.

In Fig. 1 we compare the $h^0\to\gamma\gamma$ decay width to the $\phi^0\to\gamma\gamma$ decay width, for $m_t=100$ GeV, $M=200$ GeV, $\tan\beta=1.5,4$, and the selection of μ values given in Eq. (2.2). We find, as expected from the more or less SM couplings of h^0 to WW, that the $\gamma\gamma$ width of the h^0 is dominated by the W-loop contributions for most parameter choices, and fairly similar in magnitude to that of the ϕ^0. Since the h^0 is also generally quite light new non-SM decay channels are usually absent. As a result, its $\gamma\gamma$ branching ratio can also be rather similar to that of a SM Higgs boson of the same mass. The $h^0\to\gamma\gamma$ width can, however, also be somewhat suppressed relative to that of a SM Higgs boson of similar mass due to cancellations between the W loop and the supersymmetric chargino loops. For parameter choices resulting in an extremely light chargino, the chargino loops can actually dominate the W loop and give a very substantial $h^0\to\gamma\gamma$ width. However, the mass of the lightest chargino is always so small when this is the case that it is in conflict with the experimental lower bound mentioned earlier.

In Figs. 2 and 3 we compare the $H^0, A^0\to\gamma\gamma$ decay widths to the $\phi^0\to\gamma\gamma$ decay width, for the same parameter choices as used for the h^0 comparison. In the case of the A^0, we find that the many chargino and squark loops, in combination with the other contributions, can come close to making the A^0 two-photon decay *width* as large as that found for the SM Higgs boson and the h^0. This is also true for the H^0 at large m_{H^0}, but for small m_{H^0} the $\gamma\gamma$ width is always substantially smaller than that for a ϕ^0 of the same mass. This feature at small m_{H^0} is due to a fairly complete cancellation between the mixing-angle-suppressed W loop and the top and chargino loops. At large m_{H^0} the W loop is completely negligible and this cancellation no longer occurs. Unfortunately, even when the A^0 or H^0 $\gamma\gamma$ width is comparable to that for the ϕ^0, we shall find in the following section that an adequate $\gamma\gamma$ *branching ratio* cannot be achieved for most choices of the chargino mass parameters. This is because the larger $\gamma\gamma$ widths are only achieved by choosing parameters such that the charginos (and, consequently, the neutralinos, see Ref. 4) are quite light. As a result, the Higgs-boson decays to chargino and neutralino pairs are phase-

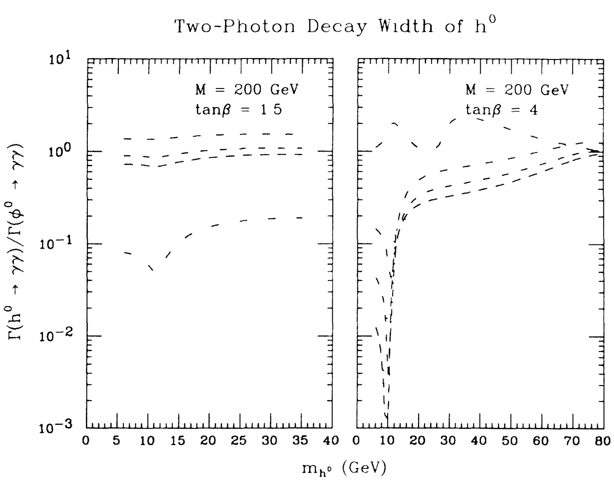

FIG. 1. The ratio $\Gamma(h^0\to\gamma\gamma)/\Gamma(\phi^0\to\gamma\gamma)$ is plotted as a function of m_{h^0} (with $m_{\phi^0}=m_{h^0}$). We take $m_t=100$ GeV, $M=200$ GeV, and give results for $\tan\beta=1.5$ and $\tan\beta=4$. In each case, five curves appear corresponding to the five μ values of Eq. (2.2). The curves are labeled in order of increasing μ by dashed, dotted, dashed-triple-dotted, dashed-double-dotted, and dashed-dotted lines, respectively.

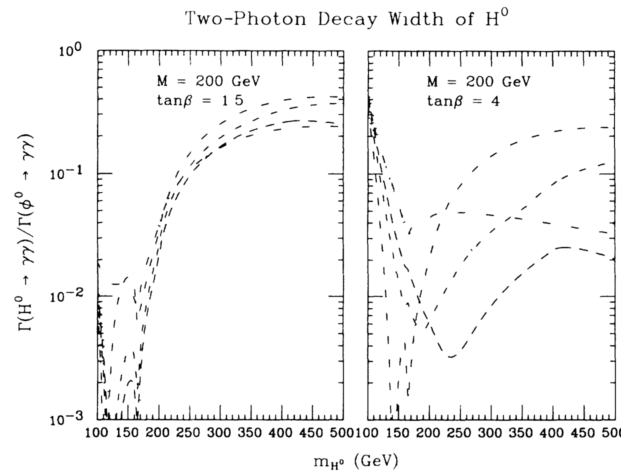

FIG. 2. The ratio $\Gamma(H^0\to\gamma\gamma)/\Gamma(\phi^0\to\gamma\gamma)$ is plotted as a function of m_{H^0} (with $m_{\phi^0}=m_{H^0}$). Parameter choices and curve labeling are the same as in Fig. 1.

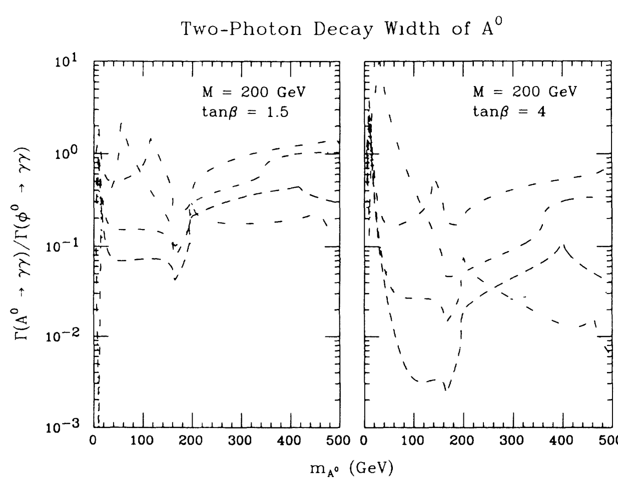

FIG. 3. The ratio $\Gamma(A^0\to\gamma\gamma)/\Gamma(\phi^0\to\gamma\gamma)$ is plotted as a function of m_{A^0} (with $m_{\phi^0}=m_{A^0}$). Parameter choices and curve labeling are the same as in Fig. 1.

space allowed. When allowed, such decays actually have larger widths than the SM heavy quark decays due to the large number of available channels and the larger effective coupling of a typical MSSM neutral Higgs boson to a chargino-chargino or neutralino-neutralino channel compared to its coupling to a heavy-quark pair [$\gtrsim g$ compared to $\gtrsim g m_q/(2m_W)$]. (For more details see Ref. 4.) The large widths associated with chargino and neutralino pair channels suppress the overall branching ratio for the $\gamma\gamma$ mode.

III. NUMERICAL RESULTS FOR THE $\gamma\gamma$ BRANCHING RATIOS

In this section we present numerical results for the $\gamma\gamma$ branching ratios of h^0, H^0, and A^0. In all cases, we have included in the calculation not only decay channels containing standard-model particles, but also decay channels containing neutralino and chargino pairs. (We have chosen the squark masses to be sufficiently heavy that squark pair channels are not allowed in the mass range of interest.) For a full discussion of these additional channels we refer to Ref. 4. As stated earlier, the loop contributions to the $\gamma\gamma$-decay widths are computed following the formalism of Ref. 5, using the parameter specifications outlined in the preceding section. All results will assume a top-quark mass of 100 GeV. In considering the sensitivity of our conclusions to this choice, it is necessary to account for the dependence of both the $\gamma\gamma$ widths and $B_{\min}$ upon m_t. We have found that quite similar conclusions to those discussed below are reached for other choices of m_t.

We begin by presenting a series of curves for $\tan\beta = 1.5$ and $m_t = 100$ GeV. These appear in three figures, Figs. 4–6. We see that for h^0, B never exceeds the minimum

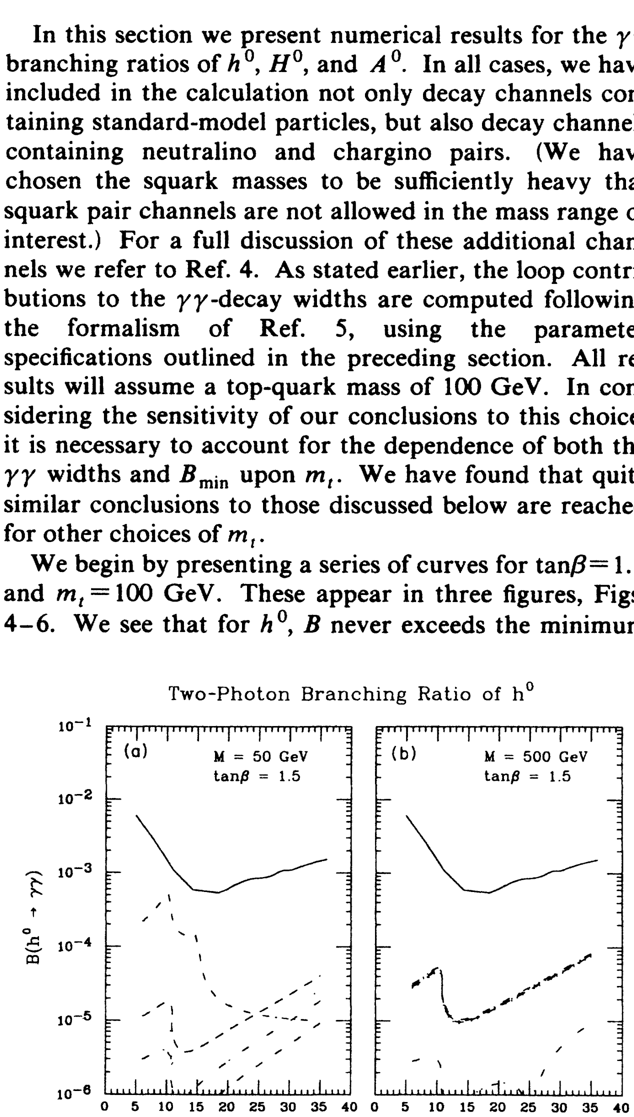

FIG. 4. We plot $B(h^0 \to \gamma\gamma)$ and $B_{\min}^{h^0}$ as a function of m_{h^0} for (a) $M = 50$ GeV and (b) $M = 500$ GeV. In each case, five curves appear corresponding to the five μ values of Eq. (2.2). The curves are labeled in order of increasing μ by dashed, dotted, dashed-triple-dotted, dashed-double-dotted, and dashed-dotted lines, respectively. We have chosen $\tan\beta = 1.5$ and $m_t = 100$ GeV. In this figure, and others to follow, the solid curve gives $B_{\min}$, see Eq. (1.5), for the Higgs boson being considered, computed using the m_t, $\tan\beta$, and $M_{\tilde{q}}$ values specified.

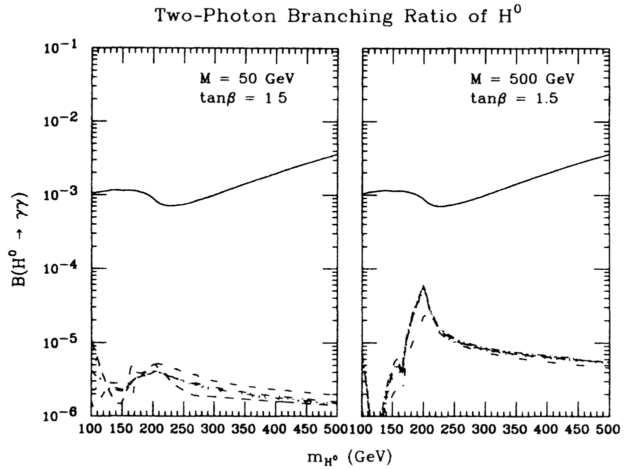

FIG. 5. We plot $B(H^0 \to \gamma\gamma)$ and $B_{\min}^{H^0}$ as a function of m_{H^0} for the same parameter selections specified in Fig. 4. See the caption of Fig. 4 for the curve legend.

required as given by the solid curve, and only approaches the needed level when the lightest chargino is below the experimental bound. For A^0, at $M = 500$ GeV $B_{\min}$ is exceeded only in a narrow region near $m_{A^0} \sim 2m_t$. Finally, for H^0, B never exceeds 10^{-5} in the case of $M = 50$ GeV and 10^{-4} in the case of $M = 500$ GeV, and is certainly always well below $B_{\min}$.

Another set of three graphs will serve to illustrate the dependence of these results on $\tan\beta$. We fix $M = 200$ GeV and compare $\tan\beta = 1.5$ to $\tan\beta = 4$ at $m_t = 100$ GeV in Figs. 7–9. Generally speaking, the $\tan\beta = 4$ branching ratios for the $\gamma\gamma$ mode are smaller than those for $\tan\beta = 1.5$ (in part, because the Higgs-boson couplings to the $t\bar{t}$ loop decrease as $1/\tan\beta$), while the required minimum branching ratios are larger (again, because of a decrease in the $t\bar{t}$ coupling which decreases the t-loop contribution to the gluon-gluon fusion cross section). For

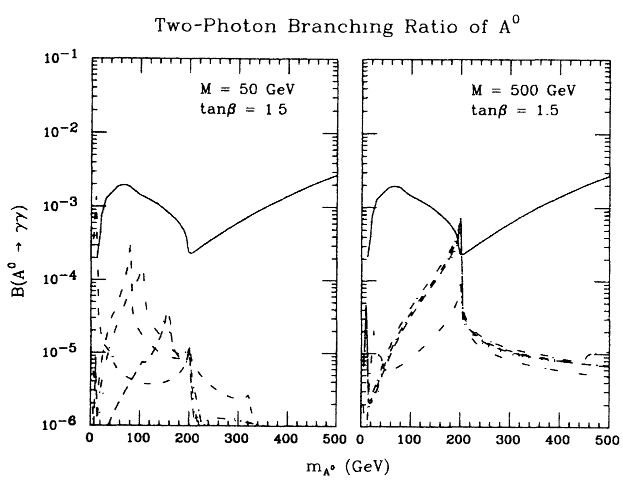

FIG. 6. We plot $B(A^0 \to \gamma\gamma)$ and $B_{\min}^{A^0}$ as a function of m_{A^0} for the same parameter selections specified in Fig. 4. See the caption of Fig. 4 for the curve legend.

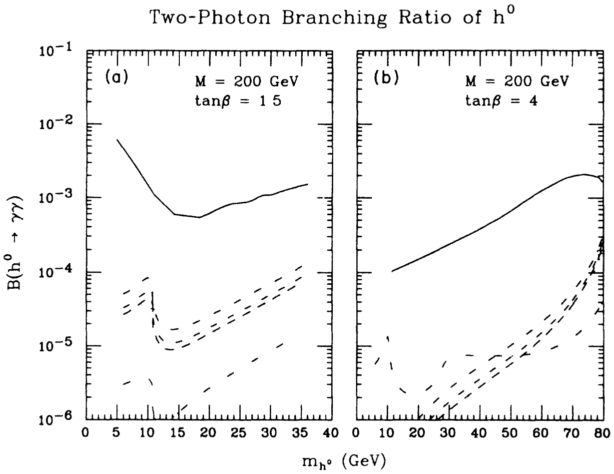

FIG. 7. We present $B(h^0 \rightarrow \gamma\gamma)$ and $B_{\min}^{h^0}$ as a function of m_{h^0} for (a) $\tan\beta = 1.5$ in comparison to (b) $\tan\beta = 4$. We have taken $M = 200$ GeV and $m_t = 100$ GeV. The labeling of curves is the same as in Fig. 4.

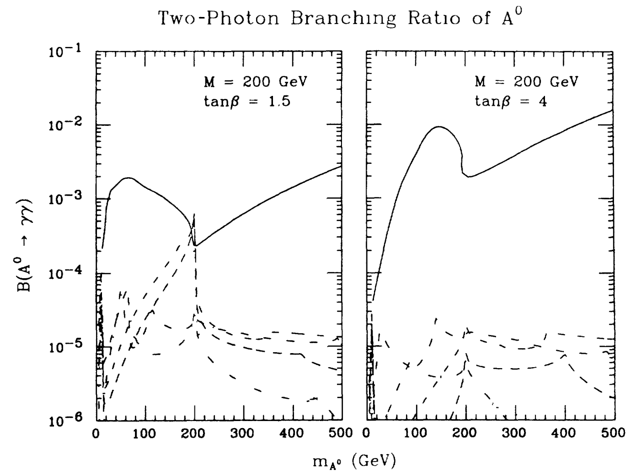

FIG. 9. We present $B(A^0 \rightarrow \gamma\gamma)$ and $B_{\min}^{A^0}$ as a function of m_{A^0} for the same parameter choices as in Fig. 7. See Fig. 4 for the labeling conventions for the curves.

this particular value of $M = 200$ GeV the $\gamma\gamma$ branching ratios are all always below $B_{\min}$ except in the case of the A^0 near $m_{A^0} \sim 2m_t$.

IV. CONCLUSIONS AND DISCUSSION

Since VV decays for the Higgs bosons appearing in the minimal supersymmetric extension of the standard model are never important or are entirely absent, it might have been that the branching ratio for decays to two photons could have been large enough to be useful for their detection, at least for masses of the neutral Higgs bosons below the $t\bar{t}$ threshold. This is certainly the case for the SM Higgs boson ϕ^0 when it has $m_{\phi^0} < \min\{2m_t, 2m_W\}$. However, generally speaking, we find that two-photon decays of the Higgs bosons of minimal supersymmetry do not provide a feasible detection mode. First, consider the h^0.

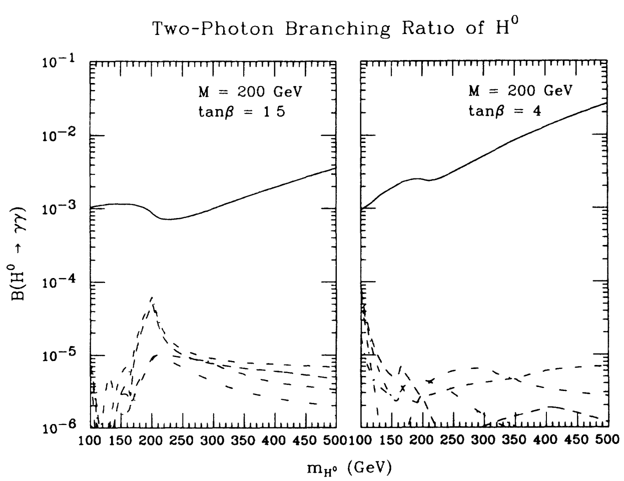

FIG. 8. We present $B(H^0 \rightarrow \gamma\gamma)$ and $B_{\min}^{H^0}$ as a function of m_{H^0} for the same parameter choices as in Fig. 7. See Fig. 4 for the labeling conventions for the curves.

The h^0 has nearly SM-like couplings to the WW loop and hence SM-like $\gamma\gamma$ width and B, at equivalent masses $m_{h^0} = m_{\phi^0}$. However, the MSSM predicts that the h^0 mass lies below m_Z. In this region of mass ($\leq m_Z$) the $\gamma\gamma$ mode cannot be used for the ϕ^0, because of the large $\gamma\gamma$ continuum background. This implies that the h^0 too cannot be discovered over the m_{h^0} range allowed by the MSSM. Turning to the H^0 and A^0, obtaining a large $\gamma\gamma$-decay width tends to require a large contribution from the chargino loops contributing to such decays, which in turn requires a light mass for the charginos. (As we have noted, non-SM loop contributions are needed to obtain large two-photon decay widths for the H^0 or A^0 because the W loop, which is the dominant loop in the case of a SM Higgs and for the h^0, is suppressed or absent, respectively.) As a result, the Higgs boson decays to a chargino and neutralino pairs become important and suppress the $\gamma\gamma$ branching ratio. Only in isolated regions of parameter space does one achieve $\gamma\gamma$ branching ratios significantly larger than 10^{-4}. As indicated by the $B_{\min}$ curves in each of the branching-ratio figures, this is not adequate to allow Higgs detection in the $\gamma\gamma$ mode. Thus, we have not pursued the $\gamma\gamma$-decay mode further.

Of course, there is always the question as to how sensitive we are to the specific structure of the minimal supersymmetric model. The three most important features of the model that influenced our calculation were (1) the weak coupling of all but the lightest scalar Higgs boson to WW and ZZ pairs, (2) the more or less standard-model strength of the $q\bar{q}$ couplings to the various neutral Higgs bosons, and (3) the importance of neutralino and chargino pair decays for the neutral Higgs bosons. A number of extensions of the MSSM have been considered, based on superstring-motivated grand unification groups. These are reviewed in Ref. 13. Generally speaking, most of the above crucial features of the MSSM model are preserved in these extended models, and we would anticipate rather similar results. In particular, in the models studied to date: (i) any Higgs boson with standard-model-like coupling to the crucial W loop is almost certain to have mass below ~ 100 GeV, where backgrounds to the $\gamma\gamma$ decay mode are large; and (ii) chargino and neutralino decay modes become a dominant decay channel when kinematically allowed, so that large chargino loop contributions to the $\gamma\gamma$ width coming from a light chargino are almost certainly coupled with large decay widths to $\tilde{\chi}\tilde{\chi}$ modes which, in turn, suppress the $\gamma\gamma$ branching ratio.

ACKNOWLEDGMENTS

One of us (J.F.G.) would like to thank the Institute for Theoretical Physics at University of California, Santa Barbara for its hospitality during the course of writing this paper. Two of us (G.G. and S.F.N.) would like to thank the Lawrence Berkeley Laboratory theoretical physics group for its kind hospitality. This work was supported in part by the Department of Energy under Grants Nos. DOE-76ER-70191-MODA33 and DE-AC03-76SF-00098, and by the National Science Foundation under Grant No. PHY82-17853, supplemented by funds from the National Aeronautics and Space Administration. G.G. was supported by the Istituto Nazionale di Fisica Nucleare (INFN, Italy). S.F.N. was supported by the Conselho Nacional de Desenvolvimento Cientifico e Tecnologico (CNPq, Brazil).

APPENDIX A: CONVENIENT EXPRESSIONS FOR THE HIGGS-BOSON WIDTHS

We write

$$\Gamma(h\to\gamma\gamma)=\frac{(\alpha M_h)^3}{(4\pi\sin\theta_W m_W)^2}\left|\sum_i I_h^i\right|^2 , \tag{A1}$$

where the sum runs over different types of loop contributions as specified below.

Contributions come from charged-Higgs-boson loops:

$$I_h^{H^\pm}=-R_{H^\pm}^h\left[\frac{m_W}{m_{H^\pm}}\right]^2\lambda[1+2\lambda F(\lambda)] ; \tag{A2}$$

sfermion loops:

$$I_h^{\rm sf}=-N_c Q_f^2 R_{sf}^h\left[\frac{m_Z}{M_{sf}}\right]^2\lambda[1+2\lambda F(\lambda)] ; \tag{A3}$$

fermion loops:

$$I_h^f=N_c Q_f^2 R_f^h[2\xi\lambda+\lambda(4\xi\lambda-1)F(\lambda)] , \tag{A4}$$

where $\xi=0$ (1) for a pseudoscalar (scalar) Higgs boson; chargino loops:

$$I_h^\chi=R_\chi^h\frac{m_W}{M_\chi}[2\xi\lambda+\lambda(4\xi\lambda-1)F(\lambda)] ; \tag{A5}$$

and W loop:

$$I_h^W=R_W^h[3\lambda(1-2\lambda)F(\lambda)-3\lambda-1/2] . \tag{A6}$$

In the above,

$$F(\lambda)=\begin{cases}-2\left[\arcsin\dfrac{1}{2\sqrt{\lambda}}\right]^2 & \text{for } \lambda\geq\frac{1}{4} ,\\ \frac{1}{2}\left[\ln\left[\dfrac{1+\sqrt{1-4\lambda}}{1-\sqrt{1-4\lambda}}\right]+i\pi\right]^2 & \text{for } \lambda<\frac{1}{4} ,\end{cases} \tag{A7}$$

with $\lambda=M^2/M_h^2$, where M is the mass of the particle running in the loop. The factors Q_f and N_c appearing in the fermion and sfermion loops are the fermion charge and number of colors, respectively. The constants R^h appearing in the various cases above are given below. For $h=h^0$:

$$R_{H^\pm}^{h^0}=-\sin(\alpha-\beta)+\frac{\cos2\beta\sin(\alpha+\beta)}{2\cos^2\theta_W} , \tag{A8}$$

$$R_{\rm sf(left)}^{h^0}=\left[\frac{M_f}{m_Z}\right]^2R_f^{h^0}-(T_f^3-Q_f\sin^2\theta_W)\sin(\alpha+\beta) , \tag{A9}$$

$$R_{\rm sf(right)}^{h^0}=\left[\frac{M_f}{m_Z}\right]^2R_f^{h^0}-Q_f\sin^2\theta_W\sin(\alpha+\beta) , \tag{A10}$$

$$R_{f(\rm up)}^{h^0}=\frac{\cos\alpha}{\sin\beta} , \tag{A11}$$

$$R_{f(\rm down)}^{h^0}=-\frac{\sin\alpha}{\cos\beta} , \tag{A12}$$

$$R_\chi^{h^0}=2(S_{ii}\cos\alpha-Q_{ii}\sin\alpha) , \tag{A13}$$

$$R_W^{h^0}=-\sin(\alpha-\beta) . \tag{A14}$$

For $h=H^0$:

$$R_{H^\pm}^{H^0}=\cos(\alpha-\beta)-\frac{\cos2\beta\cos(\alpha+\beta)}{2\cos^2\theta_W} , \tag{A15}$$

$$R_{\rm sf(left)}^{H^0}=\left[\frac{M_f}{m_Z}\right]^2R_f^{H^0}+(T_f^3-Q_f\sin^2\theta_W)\cos(\alpha+\beta) , \tag{A16}$$

$$R_{\rm sf(right)}^{H^0}=\left[\frac{M_f}{m_Z}\right]^2R_f^{H^0}+Q_f\sin^2\theta_W\cos(\alpha+\beta) , \tag{A17}$$

$$R_{f(\rm up)}^{H^0}=\frac{\sin\alpha}{\sin\beta} , \tag{A18}$$

$$R_{f(\rm down)}^{H^0}=\frac{\cos\alpha}{\cos\beta} , \tag{A19}$$

$$R_\chi^{H^0}=2(Q_{ii}\cos\alpha+S_{ii}\sin\alpha) , \tag{A20}$$

$$R_W^{H^0}=\cos(\alpha-\beta) . \tag{A21}$$

For $h=A^0$:

$$R_{H^\pm}^{A^0}=0 , \tag{A22}$$

$$R_{\rm sf(left)}^{A^0}=0 , \tag{A23}$$

$$R_{\rm sf(right)}^{A^0}=0 , \tag{A24}$$

$$R_{f(\mathrm{up})}^{A^0}=\cot\beta\ , \tag{A25}$$

$$R_{f(\mathrm{down})}^{A^0}=\tan\beta\ , \tag{A26}$$

$$R_{\chi}^{A^0}=-2(Q_{ii}\sin\beta+S_{ii}\cos\beta)\ , \tag{A27}$$

$$R_{W}^{A^0}=0\ , \tag{A28}$$

with Q_{ij}, S_{ij} and the mixing angle α as defined in Refs. 2 and 3.

[1]For a review and references, see H. P. Nilles, Phys. Rep. **110**, 1 (1984); H. E. Haber and G. L. Kane, *ibid.* **117**, 75 (1985).

[2]J. F. Gunion and H. E. Haber, Nucl. Phys. **B272**, 1 (1986).

[3]J. F. Gunion and H. E. Haber, Nucl. Phys. **B278**, 449 (1986).

[4]J. F. Gunion and H. E. Haber, Nucl. Phys. **B307**, 445 (1988).

[5]J. F. Gunion, L. Roszkowski, A. Turski, H. E. Haber, G. Gamberini, B. Kayser, S. F. Novaes, F. Olness, and J. Wudka, Phys. Rev. D **38**, 3444 (1988).

[6]J. F. Gunion, H. E. Haber, F. E. Paige, Wu-Ki Tung, and S. S. D. Willenbrock, Nucl. Phys. **B294**, 621 (1987).

[7]J. F. Gunion, G. L. Kane, and J. Wudka, Nucl. Phys. **B299**, 231 (1988).

[8]See the reports by C. Baltay, J. Huston, and B. G. Pope, in *Physics of the Superconducting Supercollider, Snowmass, 1986,* proceedings of the Summer Study, Snowmass, Colorado, 1986, edited by R. Donaldson and J. Marx (Division of Particles and Fields of the APS, New York, 1987), p. 355; D. Carlsmith *et al.*, *ibid.*, p. 405.

[9]CDF review talk by Bob Blaire, 1988 Spring Meeting of the American Physical Society, Baltimore, Maryland (unpublished).

[10]D. A. Dicus and S. S. D. Willenbrock, Phys. Rev. D **37**, 1801 (1988).

[11]R. Bates, J. N. Ng, and P. Kalyniak, Phys. Rev. D **33**, 755 (1986); **34**, 172 (1986).

[12]W. Bartel *et al.*, Z. Phys. C **29**, 505 (1985); H. J. Behrend *et al. ibid.* **35**, 181 (1985).

[13]See J. F. Gunion, Report No. NSF-ITP-88-28, 1988 (unpublished).