Detection of minimal supersymmetric model Higgs bosons in $\gamma\gamma$ collisions: Influence of supersymmetric decay modes

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We demonstrate that supersymmetric decay modes of the neutral Higgs bosons of the MSSM could well make their detection extremely difficult when produced singly in $\gamma\gamma$ collisions at a backscattered laser beam facility.

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I. INTRODUCTION

Supersymmetric (SUSY) models are leading candidates for extending the standard model (SM). The simplest such model is the minimal supersymmetric model (MSSM), which is defined by having precisely two Higgs doublets [1]. The physical Higgs eigenstates comprise two charged Higgs bosons ($H^{\pm}$), two CP-even Higgs bosons (h^0 and H^0 with $m_{h^0} < m_{H^0}$), and one CP-odd Higgs boson (A^0). A possibly very important means for discovering the neutral MSSM Higgs bosons at an e^+e^- collider is to produce them via collisions of polarized photons [2,3] obtained by backscattering polarized laser beams off of polarized electron and positron beams at a TeV-scale linear e^+e^- collider [4,5]. In previous work [2], it has been established that the neutral MSSM Higgs bosons can indeed be detected in $\gamma\gamma$ collisions over much of parameter space, *provided they decay primarily to SM final states.* In fact, since the possibly heavy H^0 and A^0 can be produced singly by direct $\gamma\gamma$ collisions, whereas they are only detectable in e^+e^- collisions in the pair production mode, $e^+e^- \to A^0H^0$, photon-photon colliders can even provide a larger discovery mass reach than direct e^+e^- collisions [2]. However, an open question is the extent to which the possibilities for A^0 and H^0 detection in $\gamma\gamma$ collisions are altered by significant decays to supersymmetric particle channels. In this paper, we show that such decays could have a decidedly negative impact.

The importance of supersymmetric decays of the MSSM Higgs bosons is dictated by the parameters of soft supersymmetry breaking. The four basic parameters are (a) the gaugino masses M_a (where a labels the group), (b) the scalar masses m_i (where i labels the various scalars, e.g., Higgs bosons, sleptons, squarks); (c) the soft Yukawa coefficients A_{ijk}, and (d) the B parameter which specifies the soft mixing term between the two Higgs scalar fields. The success of gauge coupling unification in the context of the MSSM lends considerable credence not only to the possibility that this extension of the standard model is correct, but also to the idea that the boundary conditions for all the soft-supersymmetry-breaking parameters at the unification scale could be relatively simple and universal. Superstring theory provides particularly attractive and well-motivated examples of such boundary conditions. In this paper we consider the dilatonlike superstring supersymmetry-breaking scheme (labeled by D). This is one of the most attractive models available and yields a complex array of decay channels for the MSSM Higgs bosons. In this model the M_a, m_i, and A_{ijk} parameters all take on universal values at the unification scale M_U related by

$$M^0 = -A^0 = \sqrt{3}m^0. \quad (1)$$

Predictions in this model for the B parameter are rather uncertain, and so it is kept a free parameter. The dilatonlike boundary conditions are certainly those appropriate when supersymmetry breaking is dominated by the dilaton field in string theory, but they also apply for a remarkably broad class of models (including Calabi-Yau compactifications, and orbifold models in which the MSSM fields all belong to the untwisted sector), so long as the moduli fields do not play a dominant role in supersymmetry breaking. For a brief review and detailed references, see Ref. [6].

If the boundary conditions of Eq. (1) are imposed and the top quark mass is fixed [we adopt $m_t(m_t) = 170$ GeV, corresponding to a pole mass of about 178 GeV], only two free parameters and a sign remain undetermined after minimizing the potential. The two parameters can be taken to be $\tan\beta$, the ratio of the neutral Higgs field vacuum expectation values, and $m_{\tilde{g}}$, the gluino mass. The parameter B is determined in terms of these, as are all other superpotential parameters, including the magnitude of the Higgs superfield mixing parameter μ. However, the sign of μ is not determined. Two models results, D^+ and D^-, the superscript indicating the sign of μ, the phenomenology of which can be explored in the two-dimensional $m_{\tilde{g}}$-$\tan\beta$ parameter space.

The discussion so far has obscured one fundamental problem facing the gauge coupling unification success: namely, the scale M_U at which the couplings naturally unify is $\sim 2 \times 10^{16}$ GeV, i.e., much less than the natural scale for supergravity and string unification of $M_S \sim 10^{18}$ GeV. A variety of excuses for this have been discussed. In Ref. [6] two extreme approaches were adopted: (i) ignore the difference—a more complete understanding of the feed-down of SUSY breaking from the full supergravity or superstring theory could resolve the discrepancy;

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(ii) assume that the unification at M_U is only apparent (i.e., accidental) and introduce a minimal set of additional matter fields at high scale with masses chosen precisely so as to give coupling unification at M_S. We will not go into detail regarding these extra fields; a discussion and references can be found in Ref. [6]. The models with such extra fields are termed the "string-scale-unified" versions of the previously listed models, and will be denoted by SD^+ and SD^-.

To systematically investigate the resulting models, Ref. [6] first established the allowed region of $m_{\tilde{g}}$-tanβ parameter space for each subject to (a) all predicted SUSY partner particles (including the light Higgs boson h^0) are unobservable, (b) the lightest SUSY particle is either the lightest neutralino $\tilde{\chi}_1^0$ (as is always the case for the allowed parameter space of the models explored here) or the sneutrino $\tilde{\nu}$, (c) the top quark Yukawa coupling remains perturbative at all scales from m_W up to M_U or M_S, and (d) proper electroweak symmetry breaking and a global minimum are obtained. Constraints from $b \to s\gamma$, relic abundance, and proton decay were not imposed, as these all have considerable uncertainties and/or require additional model-dependent input. Exact $b - \tau$ Yukawa unification was also not imposed.

Within the allowed parameter spaces, the masses of the SUSY particles scale with $m_{\tilde{g}}$; variation of the masses with tanβ at fixed $m_{\tilde{g}}$ is relatively limited, especially for $m_{\tilde{g}}$ values above about 500–600 GeV, with $\tilde{l}_R, \tilde{\chi}_1^+, \tilde{\chi}_2^0, \tilde{\nu}, \tilde{l}_L$ clustering between 0.2 to 0.4 times $m_{\tilde{g}}$. It is the restricted size of the soft scalar mass parameter m^0 relative to M^0 that causes the sleptons to be rather light in the dilatonlike models. Indeed, slepton masses are largely generated by renormalization-group evolution from the M^0 gaugino seed value at M_U; only the squarks acquire masses comparable to $m_{\tilde{g}}$, as a result of the driving terms proportional to α_s in the renormalization-group equations (RGE's).

Regarding the Higgs boson masses, a very general pattern emerges. The h^0 is normally relatively light, even after including the standard one-loop radiative corrections [7], which depend most crucially upon the top quark mass (m_t) and the stop squark mass ($m_{\tilde{t}}$). For gluino masses below 1 TeV and m_t (m_t) = 170 GeV, $m_{h^0} \leq 125$ GeV, with quite low values ($65 \lesssim m_{h^0} \lesssim 110$ GeV) being rather typical. Thus, the h^0 will be easily discovered via $e^+e^- \to Zh^0$ (even if the h^0 decays invisibly to $\tilde{\chi}_1^0\tilde{\chi}_1^0$, as can happen in these models). In contrast, the RGE driven electroweak symmetry-breaking models in general, and the dilatonlike boundary condition models in particular, predict rather large $m_{A^0} \sim m_{H^0} \sim m_{H^\pm}$ values. For most of parameter space, $m_{A^0} \gtrsim 200$ GeV with values in the 300–600 GeV range being much more typical for $m_{\tilde{g}} < 800$ GeV.[1] This means that $e^+e^- \to A^0H^0, H^+H^-$ pair production is quite possibly disallowed kinematically for a $\sqrt{s} \sim 500$ GeV e^+e^- collider, and that single production via $\gamma\gamma \to A^0, H^0$ would be the only possible mode of discovery. Further, for $m_{\tilde{g}} \lesssim 800$ GeV, the $\tilde{\chi}_1^0, \tilde{\chi}_2^0, \tilde{\chi}_1^\pm, \tilde{\nu}, \tilde{l}_R$, and (except at high tanβ) $\tilde{l}_L$ are all light enough to appear in two-body decay modes of the A^0 and H^0. Thus, the D and SD models present many possible scenarios of precisely the type that we wish to explore.

[1]This upper bound represents a purely aesthetical choice as to an $m_{\tilde{g}}$ value below which the model is clearly not fine tuned.

II. SCENARIOS

Of the specific $m_{\tilde{g}}$-tanβ scenarios explored with regard to their general phenomenology in Ref. [6], we focus on a limited number of representative cases. In the notation of Ref. [6], these are the scenarios D_3^+, D_1^-, D_4^-, D_4^+, D_5^+, D_7^+, SD_1^-, and SD_2^-, where we have listed them in order of increasing m_{A^0}. A complete listing of all relevant particle masses, and a summary of the decay modes of the SUSY particles is presented in Ref. [6]. Here, we give a condensed summary along with details regarding the decays of the A^0 and H^0 Higgs bosons. The scenarios are summarized in Table I, where we give masses for the Higgs bosons and selected superparticles.

Detailed decay tables for the Higgs bosons and superparticles were generated using ISASUSY [8], and cross checked using expanded versions of the programs developed for the work of Ref. [2]. The important Higgs branching ratios as a function of scenario are presented in Table II. Note that the cumulative effect of the SUSY decay modes is generally to substantially reduce the SM particle modes, unless tanβ is very large (as in the D_3^+ and SD_1^- cases), in which case the $b\bar{b}$ mode can still be dominant. Especially dramatic is the dominance of the $\tilde{\nu}\tilde{\nu}$ decay modes in the D_1^- and D_4^- cases, which has a drastic impact given that in these cases the $\tilde{\nu}$ itself decays invisibly (see Ref. [6]).

The formalism for computing the rate of Higgs boson production in $\gamma\gamma$ collisions is well-established [1,2]. An approximate result[2] for the number of Higgs bosons produced at a backscattered laser-beam facility is

$$N(\gamma\gamma \to h) = \frac{4\pi^2\Gamma(h \to \gamma\gamma)}{m_h^3} y_h F(y_h)(1 + \langle\lambda\lambda'\rangle_{y_h})L_{e^+e^-}, \tag{2}$$

where $y_h \equiv m_h/E_{e^+e^-}$, and $F(y_h)$ and $\langle\lambda\lambda'\rangle_{y_h}$ are obtained by convoluting together the spectra and polarizations for the backscattered photons. In computing $\Gamma(h \to \gamma\gamma)$, the full set of SUSY and SM particle loops is included. For each given scenario these contributions are completely known, since all parameters and masses of the MSSM are fixed. In computing $F(y_h)$ we have been as optimistic as possible, choosing the laser-photon polarizations, e^+ and e^- polarizations, and machine energy so that the $\gamma\gamma$ spectrum is sharply peaked and is centered at the Higgs boson mass of interest. The most highly peaked spectrum is obtained by choosing large polarizations for the e^+ and e^- (we adopt $\lambda_e = \lambda_e' = +0.45$),

[2]In practice, we employ a more accurate numerical procedure.

TABLE I. A tabulation of supersymmetric particle masses (in GeV) for the D and SD scenarios considered.

Scenario	$m_{\tilde{g}}$	$\tan\beta$	m_{h^0}	m_{A^0}	m_{H^0}	$m_{\tilde{\chi}_1^0}$	$m_{\tilde{\chi}_2^0}$	$m_{\tilde{\chi}_1^+}$	$m_{\tilde{l}_L}$	$m_{\tilde{l}_R}$	$m_{\tilde{\nu}}$	$m_{\tilde{q}}$	$m_{\tilde{t}_1}$
D_3^+	310	15.0	103	180	180	39.9	72.5	70.2	109	85.9	74.4	277	188
D_1^-	232	2.0	58.4	190	205	37.1	83.5	83.3	82.3	65.0	54.1	207	215
D_4^-	301	2.2	69.0	244	255	47.3	100	100	103	80.2	79.8	269	242
D_4^+	346	3.2	93.6	250	255	40.4	79.2	73.5	118	91.8	93.0	310	195
D_5^+	431	4.5	104	300	302	58.4	109	107	144	111	122	386	250
D_7^+	503	5.0	108	350	351	71.3	134	133	166	127	147	450	297
SD_1^-	471	15.0	111	357	357	69.1	134	134	193	157	176	464	301
SD_2^-	503	5.0	105	424	426	75.4	149	149	205	166	190	496	339

large polarizations (opposite those for the e^+, e^-) for the laser photons (we taken $P_c = P_c' = -1$), and as large a value for the ξ parameter as possible (we employ $\xi = 4.8$) without going above pair production threshold.[3] (For details see Refs. [3–5].) For these choices, the spectrum is peaked in the vicinity of $y_h = 0.79$ for which $y_h F(y_h)(1 + \langle\lambda\lambda'\rangle_{y_h}) \sim 3.5$, with $\langle\lambda\lambda'\rangle_{y_h} \sim 0.94$. [The corresponding value of $F\ (y_h) \sim 2.3$ is illustrated, for example, in Fig. 9(d) of Ref. [3], for a very similar backscattered laser-beam configuration.]

The resulting total rates for A^0 and H^0 production for each scenario appear in Table III (assuming an integrated luminosity of $L \equiv L_{e^+e^-} = 10$ fb^{-1}, such as might be accumulated in one year of operation), along with the corresponding choices of optimal $\sqrt{s}$ for the e^+e^- collider. Note that the decline in production rate with increasing Higgs boson mass due to the m_h^{-3} factor in Eq. (2) is significantly modulated by variations in $\Gamma(h \to \gamma\gamma)$, which in particular is sharply suppressed at large $\tan\beta$ due to enhanced cancellations from the b-quark loop contribution, whereas it turns out to be comparatively enhanced for the SD_2^- scenario.

We recognize that the use of a highly-peaked spectrum for initial discovery of the Higgs bosons is unrealistic in practice, as it requires scanning in order to discover a given Higgs boson. However, we have adopted a highly-peaked spectrum for two reasons. First, it yields the most optimistic results possible, which will not prove to be terribly promising. Second, it gives an accurate representation of what would be possible should the mass of a given Higgs boson already be known, in which case $\gamma\gamma$ collision detection would be a second generation experiment motivated by the importance of determining $\Gamma(h \to \gamma\gamma)$. In practice, A^0 and H^0 Higgs boson *searches* in $\gamma\gamma$ collisions (i.e., prior to their discovery elsewhere) would probably employ a fixed $\sqrt{s}$, in which case it is probably most reasonable to assume that m_{A^0} and m_{H^0} would not be $\sim 0.79\sqrt{s}$. The above-specified backscattered laser-beam configuration [for which $F(y_h)$ falls to ~ 1 for $y_h \lesssim 0.6$] would be employed in order to explore for Higgs bosons with $m_h \sim 0.6 - 0.8\sqrt{s}$, while the configuration $\lambda_e \sim \lambda_e' \sim 0.45, P_c \sim P_c' \sim +1$ [for which $F(y_h)$ exhibits a spectrum that is broadly peaked with $F(y_h) \sim 1.7$ in the vicinity of $y_h \sim 0.4$ falling below 1 for y_h below 0.1 and above 0.6—see Fig. 9(b) of Ref. [3]] would be employed to explore for Higgs bosons below $0.6\sqrt{s}$. Then, the true rates for the various channels considered here would most typically be between 20% and 50% lower than those quoted below assuming we sum over two runs with an integrated luminosity of $L = 10$ fb^{-1} in *each* of the two complementary backscattered laser-beam configurations outlined above.

We turn next to rates in specific channels. Tree-level backgrounds are present for the $b\bar{b}$, $t\bar{t}$, $\tilde{\chi}_1^+\tilde{\chi}_1^-$, and $\tilde{l}\tilde{l}$ channels. The $\tilde{\chi}_1^0\tilde{\chi}_1^0$ channel is invisible, while the $\tilde{\chi}_1^0\tilde{\chi}_2^0$ and $\tilde{\chi}_2^0\tilde{\chi}_2^0$ backgrounds only arise at one loop. The h^0h^0 and h^0Z channels we regard as background free, assuming that the h^0 and Z masses can be reconstructed with reasonable accuracy in the $b\bar{b}b\bar{b}$ and $b\bar{b}Z$ (with Z visible) modes.

We examine first the $b\bar{b}$ and $t\bar{t}$ final state decay modes and their backgrounds. The rates are summarized in Table IV. In obtaining these rates we have not included the efficiency penalty that will inevitably arise in experimentally isolating the b and t final states. Further, in estimating background rates, we have assumed a 10 GeV mass resolution, which might be achievable for $b\bar{b}$ final states but is certainly far too optimistic for the $t\bar{t}$ channel. Even with these optimistic procedures, discovery of the H^0 and A^0 appears quite difficult. The statistical significance, $N_{SD} \equiv S/\sqrt{B}$, achieved by combining the A^0 and H^0 signals (not really allowed in cases where the $\sqrt{s}$ values needed to achieve the optimal rates are somewhat different) and using the average of the two backgrounds is always below $N_{SD} = 3$, and declines to no more than $N_{SD} = 1$ or 2 at higher Higgs boson masses. Thus, even for our optimal $\gamma\gamma$ spectrum and resolution choices, roughly $L \gtrsim 60$ fb^{-1} would be required for these channels to provide viable signals for most scenarios.

Let us next examine the $A^0 \to h^0Z$ and $H^0 \to h^0h^0$ channels. Raw event rates are presented in Table V. We see immediately that these channels only show a reasonable level of promise in the case of the D_1^- and D_4^- scenarios. These two scenarios illustrate more generally the ingredients required in order that the h^0Z and h^0h^0 channels yield viable discovery signals: (i) the A^0 and H^0 masses are sufficiently modest that the m_h^{-3} factor in Eq.

[3] We remind the reader that these choices also maximize $1 + \langle\lambda\lambda'\rangle_{y_h}$, which not only enhances the Higgs boson production rate, but also minimizes all of the two-body continuum background channels of interest: $b\bar{b}, t\bar{t}, \tilde{\chi}_1^+\tilde{\chi}_1^-$, and $\tilde{l}\tilde{l}$.

TABLE II. A tabulation of important branching ratios for (a) the H^0 and (b) the A^0. In the results $\tilde{l} = \tilde{e}, \tilde{\mu}$ are summed together and all $\tilde{\nu}\tilde{\nu}$ modes are summed together. We use the shorthand notation $\tilde{l}\tilde{l} \equiv \tilde{l}_L\tilde{l}_L + \tilde{l}_R\tilde{l}_R$.

(a) H^0 branching ratios

Scenario	$b\bar{b}$	$t\bar{t}$	$W^+W^- + ZZ$	$\tilde{\chi}_1^0\tilde{\chi}_1^0$	$\tilde{\chi}_1^0\tilde{\chi}_2^0$	$\tilde{\chi}_2^0\tilde{\chi}_2^0$	$\tilde{\chi}_1^+\tilde{\chi}_1^-$	h^0h^0	$\tilde{l}\tilde{l}$	$\tilde{\nu}\tilde{\nu}$
D_3^+	0.782		0.0003	0.031	0.046	0.011	0.072		0.0003	0.003
D_1^-	0.045		0.038	0.002	0.031	0.088	0.112	0.103	0.110	0.414
D_4^-	0.072		0.054	0.004	0.053	0.105	0.155	0.149	0.081	0.280
D_4^+	0.144		0.038	0.104	0.126	0.034	0.292	0.064	0.034	0.136
D_5^+	0.343		0.024	0.062	0.136	0.060	0.247	0.028	0.014	0.054
D_7^+	0.456	0.030	0.018	0.040	0.113	0.058	0.187	0.016	0.009	0.032
SD_1^-	0.833	0.001	0.002	0.005	0.022	0.019	0.042	0.013	0.0001	0.0003
SD_2^-	0.315	0.273	0.018	0.010	0.057	0.068	0.134	0.082	0.004	0.013

(b) A^0 branching ratios

Scenario	$b\bar{b}$	$t\bar{t}$	$\tilde{\chi}_1^0\tilde{\chi}_1^0$	$\tilde{\chi}_1^0\tilde{\chi}_2^0$	$\tilde{\chi}_2^0\tilde{\chi}_2^0$	$\tilde{\chi}_1^+\tilde{\chi}_1^-$	h^0Z
D_3^+	0.726		0.040	0.076	0.034	0.075	
D_1^-	0.113		0.009	0.144	0.504	0.189	0.031
D_4^-	0.128		0.015	0.160	0.407	0.231	0.048
D_4^+	0.096		0.152	0.230	0.087	0.419	0.010
D_5^+	0.240		0.076	0.218	0.153	0.286	0.008
D_7^+	0.271	0.198	0.041	0.152	0.136	0.176	0.006
SD_1^-	0.819	0.009	0.005	0.028	0.038	0.037	0.001
SD_2^-	0.255	0.470	0.009	0.056	0.089	0.091	0.009

TABLE III. A tabulation of inclusive Higgs boson production rates as a function of scenario. We assume $L = 10\ \mathrm{fb}^{-1}$ and have optimized the $\gamma\gamma$ energy spectrum and collider energy. The corresponding optimal e^+e^- energy (in GeV) for each Higgs boson is tabulated.

Scenario	m_{A^0}	A^0 rate	$\sqrt{s}_{\rm opt}$	m_{H^0}	H^0 rate	$\sqrt{s}_{\rm opt}$
D_3^+	180	56	228	180	40	228
D_1^-	190	363	240	205	466	260
D_4^-	244	210	309	255	190	323
D_4^+	250	70	316	255	46	324
D_5^+	300	14	381	302	50	382
D_7^+	350	6	443	351	59	445
SD_1^-	357	0.5	451	357	11	452
SD_2^-	424	38	538	426	17	538

TABLE IV. A tabulation of Higgs boson signal and background rates (assuming $L = 10\ \mathrm{fb}^{-1}$) for the $b\bar{b}$ and $t\bar{t}$ channels. In computing the background rates a final state mass resolution of 10 GeV is assumed.

Scenario	$A^0 \to b\bar{b}$	Background	$H^0 \to b\bar{b}$	Background	$A^0 \to t\bar{t}$	Background	$H^0 \to t\bar{t}$	Background
D_3^+	41	770	31	770				
D_1^-	41	670	21	570				
D_4^-	27	320	14	290				
D_4^+	7	300	7	290				
D_5^+	3	180	17	170				
D_7^+	2	120	27	120	1	350	2	370
SD_1^-	0.4	110	9	110	0.005	430	0.02	430
SD_2^-	10	70	5	77	18	580	5	570

TABLE V. A tabulation of signal rates (assuming $L = 10$ fb^{-1}) in the $H^0 \to h^0 h^0$ and $A^0 \to h^0 Z$ channels.

Scenario	$A^0 \to h^0 Z$ rate	$H^0 \to h^0 h^0$ rate
D_3^+	0	0
D_1^-	11	48
D_4^-	10	28
D_4^+	0.7	2.9
D_5^+	0.1	1.4
D_7^+	0.04	0.9
SD_1^-	0.0005	0.14
SD_2^-	0.35	1.4

(2) does not yield too much rate suppression, but sufficiently large that h^0Z and h^0h^0 decays are kinematically allowed, (ii) the value of $\tan\beta$ is moderate so that the $b\bar{b}$ decay channel of the Higgs bosons does not overwhelm all others and the b-quark loop is not enhanced so as to cause cancellations that yield small values for $\Gamma(A^0, H^0 \to \gamma\gamma)$, and (iii) the Higgs boson masses are small enough that SUSY decay modes still suffer some kinematical suppression. Of course, in realistically assessing the visibility of the h^0Z and h^0h^0 signals one must take into account the fact that $h^0h^0 \to b\bar{b}b\bar{b}$ and $h^0Z \to b\bar{b}$+visible branching fractions [typically $B(h^0 \to b\bar{b}) \sim 0.9$ and $B(Z \to \text{visible}) \sim 0.8$] will reduce the effective rates for useful channels and the fact that to isolate these channels from QCD backgrounds it will be necessary to tag at least one of the b-quark jets (with roughly 60% efficiency). Consequently, the effective rates for these promising channels will be somewhat marginal even in the most favorable scenarios, unless $L > 10$ fb^{-1} is accumulated.

Could SUSY decay channels save the day? Let us first focus on the tree-level background-free $\tilde{\chi}_1^0\tilde{\chi}_2^0$ and $\tilde{\chi}_2^0\tilde{\chi}_2^0$ channels. The rates for these channels for the A^0 and H^0 are given in Table VI. In order to assess the possible utility of these rates we need to include the $\tilde{\chi}_2^0$ decays. The primary decays of the $\tilde{\chi}_2^0$ are of three basic types: $ll + E_T^{\text{miss}}$ (often via the two-body $l\tilde{l}_R$ mode, with $\tilde{l}_R \to l\tilde{\chi}_1^0$), $jj + E_T^{\text{miss}}$ (in which we include $\tau\tau + E_T^{\text{miss}}$, aside from which it is always a three-body decay), and pure E_T^{miss} (often via two-body $\tilde{\nu}\nu$ modes where the $\tilde{\nu}$ decays invisibly via $\tilde{\nu} \to \nu\tilde{\chi}_1^0$). The branching ratios for these three basic types of $\tilde{\chi}_2^0$ decay are given in Table VII as a function of scenario.

The types of Higgs boson final states that result are of six basic classes. The $\tilde{\chi}_1^0\tilde{\chi}_2^0$ decay mode of the A^0 and H^0 can lead to a purely invisible decay channel, which we discard as unusable, a channel with two leptons and missing energy, $ll + E_T^{\text{miss}}$ (where both l's come from the $\tilde{\chi}_2^0$), and a channel with two jets and missing energy, $jj + E_T^{\text{miss}}$ (where we include τ leptons in the j). The $\tilde{\chi}_2^0\tilde{\chi}_2^0$ decay mode can lead to these same final states and, in addition, a two-lepton-two-jet plus missing energy final state, $ll + jj + E_T^{\text{miss}}$, a four-lepton plus missing energy final state, $ll + ll + E_T^{\text{miss}}$, and a four-jet plus missing energy final state, $jj + jj + E_T^{\text{miss}}$. In computing the rates for these final states we combine the events coming from the A^0 and H^0—these have similar mass, and mass reconstruction in the final state is not possible due to the missing energy content. The resulting event rates for each class of final state are displayed in Table VIII.

We see that only the $ll + E_T^{\text{miss}}$ and $jj + E_T^{\text{miss}}$ channels have a non-negligible number of events, and that even these rates are very modest. The reasons for this are several, and can be traced from Tables VI and VII. For the D_1^- and D_4^- scenarios, Higgs boson production rates were high, but decays for the $\tilde{\chi}_2^0$ are completely dominated by totally invisible channels. For the other scenarios, visible $\tilde{\chi}_2^0$ decays have a substantial branching fraction but Higgs boson production rates are much more modest. We cannot say if this conspiracy is a general phenomenon, or simply specific to the dilatonlike boundary conditions employed here.

Are the $ll + E_T^{\text{miss}}$ and/or $jj + E_T^{\text{miss}}$ events sufficiently unique to provide a viable signal? We are pessimistic in this regard, since many large rate processes can potentially yield backgrounds. Consider first the $ll + E_T^{\text{miss}}$ channel. We shall see that tree-level $\tilde{\chi}_1^+\tilde{\chi}_1^-$ continuum production has a *very* high rate, and since the $\tilde{\chi}_1^\pm$ have a significant branching fraction to $l + E_T^{\text{miss}}$, we will have a large number of $ll + E_T^{\text{miss}}$ final states from this source. Even though the two leptons of a signal event both derive from a single $\tilde{\chi}_2^0$, they will not tend to be terribly well collimated due to the large role played by the E_T^{miss} component of a given $\tilde{\chi}_2^0$ decay. Thus, we believe (but we have not performed a Monte Carlo study) that event topology will not allow a sufficiently efficient means of discriminating the signal of interest from this very large background. In addition, $\tilde{l}\tilde{l}$ production also has a very high rate and also contributes to the $ll + E_T^{\text{miss}}$ channel. Regarding the $jj + E_T^{\text{miss}}$ channel, once again $\tilde{\chi}_1^+\tilde{\chi}_1^-$ will yield a background when one $\tilde{\chi}_1^\pm$ decays hadronically to

TABLE VI. A tabulation of signal rates (assuming $L = 10$ fb^{-1}) in the $\tilde{\chi}_1^0\tilde{\chi}_2^0$ and $\tilde{\chi}_2^0\tilde{\chi}_2^0$ final states, before including $\tilde{\chi}_2^0$ decay branching fractions.

Scenario	$A^0 \to \tilde{\chi}_1^0\tilde{\chi}_2^0$	$A^0 \to \tilde{\chi}_2^0\tilde{\chi}_2^0$	$H^0 \to \tilde{\chi}_1^0\tilde{\chi}_2^0$	$H^0 \to \tilde{\chi}_2^0\tilde{\chi}_2^0$
D_3^+	4	2	2	0.4
D_1^-	52	183	14	41
D_4^-	34	86	10	20
D_4^+	16	6	6	2
D_5^+	3	2	7	3
D_7^+	1	1	7	3
SD_1^-	0.01	0.02	0.2	0.2
SD_2^-	2	3	1	1

TABLE VII. A tabulation of branching ratios (B) for the three basic $\tilde{\chi}_2^0$ decay channels.

Scenario	$B\ (ll + E_T^{\rm miss})$	$B\ (jj + E_T^{\rm miss})$	$B\ (E_T^{\rm miss})$
D_3^+	0.082	0.067	0.851
D_1^-	0.017	0.006	0.977
D_4^-	0.027	0.014	0.959
D_4^+	0.301	0.187	0.510
D_5^+	0.314	0.205	0.481
D_7^+	0.266	0.204	0.530
SD_1^-	0.206	0.362	0.432
SD_2^-	0.251	0.355	0.394

two jets plus missing energy and the other decays leptonically and the lepton is "missed." In addition, $\gamma\gamma \to$ jet + jet rates are very high and will inevitably have a significant detector-dependent missing energy tail. SUSY production processes can also contribute backgrounds; for example, $\gamma\gamma \to \tilde{q}\tilde{q}$ contributes when both squarks decay to $q\tilde{\chi}_1^0$. Thus, even before inclusion of detection efficiencies, we are relatively certain that the low Higgs boson signal event rates would not constitute viable signals. (Detailed studies will not be pursued here.) Models with very different boundary conditions could perhaps yield more viable Higgs boson signal rates in these channels.

The remaining SUSY-channel possibilities are the $\tilde{\chi}_1^+\tilde{\chi}_1^-$ and $\tilde{l}\tilde{l}$ channels. Generally speaking, both primarily yield $ll + E_T^{\rm miss}$ final states (although the $\tilde{\chi}_1^+$ can decay also to jets, this mode is generally smaller than the leptonic mode). So in some sense these channels should be considered together and also combined (to the extent that the topologies do not differ much) with the $ll + E_T^{\rm miss}$ events deriving from the $\tilde{\chi}_1^0\tilde{\chi}_2^0$ and $\tilde{\chi}_2^0\tilde{\chi}_2^0$ decay channels. (Of course, in the latter case the two l's must be of the same type, whereas for the $\tilde{\chi}_1^+\tilde{\chi}_1^-$ modes they can be of different types.) For purposes of discussion, we shall keep all these different channels separate. The event rates for these channels are given in Table IX, along with the direct tree-level backgrounds, assuming a final state mass resolution of 10 GeV. Such a small resolution is undoubtedly highly unrealistic given that the $\tilde{\chi}_1^+\tilde{\chi}_1^-$ and $\tilde{l}\tilde{l}$ final states contain significant missing energy. A cursory survey of the numbers reveals the impossibility of overcoming the backgrounds. (A number of distributions for final leptons were examined to see if any dramatic increases of S/B could be achieved by appropriate cuts, but no effective cuts were found.) Even if we ignore all topology differences and add in the $\tilde{\chi}_1^0\tilde{\chi}_2^0$ and $\tilde{\chi}_2^0\tilde{\chi}_2^0$ events of the $ll + E_T^{\rm miss}$ type, the signal rates remain very small compared to the backgrounds.

III. CONCLUSIONS

We are forced to conclude that detection of the H^0 and A^0 in $\gamma\gamma$ collisions at a backscattered laser beam facility could prove extremely difficult in models where SUSY decays of the Higgs bosons are significant, unless integrated luminosities much higher than $L = 10$ fb^{-1} could be provided. For the models explored here we found that, even for a completely optimized $\gamma\gamma$ energy spectrum, for $L = 10$ fb^{-1} the $b\bar{b}$ and $t\bar{t}$ channel rates are generally reduced to too low a level relative to the corresponding continuum backgrounds to provide a viable Higgs boson signal. The SUSY decay modes themselves do not appear to have large enough rates relative to expected backgrounds. The only channels that have a significant chance of revealing a signal are the (background-free) $A^0 \to h^0 Z \to b\bar{b}Z_{\rm vis}$ and $H^0 \to h^0 h^0 \to b\bar{b}b\bar{b}$ modes, and even the most promising specific scenarios that we have examined yield only very modest event rates despite the optimization of the $\gamma\gamma$ spectrum. Considering all possible channels, for most of the scenarios examined here $L \gtrsim 50$ fb^{-1} would be needed in order to obtain at least one viable signal.

The basic problem is that once SUSY decay modes are allowed, the large number of decay channels means that no single decay channel is likely to be dominant (with the exception of the largely or completely invisible $\tilde{\nu}\tilde{\nu}$

TABLE VIII. A tabulation of rates (assuming $L = 10$ fb^{-1}) for the five classes of visible $\tilde{\chi}_1^0\tilde{\chi}_2^0 + \tilde{\chi}_2^0\tilde{\chi}_2^0$ final state after combining A^0 and H^0 production.

Scenario	$ll + E_T^{\rm miss}$	$jj + E_T^{\rm miss}$	$ll + jj + E_T^{\rm miss}$	$ll + ll + E_T^{\rm miss}$	$jj + jj + E_T^{\rm miss}$
D_3^+	0.8	0.7	0.03	0.02	0.01
D_1^-	8	3	0.05	0.06	0.008
D_4^-	7	3	0.08	0.08	0.02
D_4^+	9	6	0.9	0.7	0.3
D_5^+	5	3	0.6	0.5	0.2
D_7^+	3	3	0.4	0.3	0.2
SD_1^-	0.1	0.1	0.03	0.009	0.03
SD_2^-	2	2	0.7	0.3	0.5

TABLE IX. A tabulation of signal rates (assuming $L = 10\ \text{fb}^{-1}$) in the $\tilde{\chi}_1^+\tilde{\chi}_1^-$ and $\tilde{l}\tilde{l}$ final states. Backgrounds in these channels are also given for the (unrealistically small) final state mass resolution of 10 GeV.

Scenario	$A^0 \to \tilde{\chi}_1^+\tilde{\chi}_1^-$	Background	$H^0 \to \tilde{\chi}_1^+\tilde{\chi}_1^-$	Background	$H^0 \to \tilde{l}\tilde{l}$	Background
D_3^+	4	13000	3	13000	0.01	4400
D_1^-	69	7900	52	8100	51	6800
D_4^-	49	4800	29	4600	15	3600
D_4^+	29	6600	13	6400	2	3800
D_5^+	4	3300	12	3300	0.7	2200
D_7^+	1	1900	11	1900	0.6	1400
SD_1^-	0.02	1800	0.5	1800	0.001	780
SD_2^-	3	1200	2	1200	0.07	780

channel). Consequently, no single final state mode obtains a high event rate. The only exception to this rule arises if tanβ is large, in which case the $b\bar{b}$ decay mode is dominant for both the H^0 and A^0, and the only issue is the absolute production rate of the Higgs bosons themselves. Unfortunately, as noted previously in Ref. [2], for Higgs boson masses in the 200–500 GeV range there is a general tendency for the enhanced b-quark loop to significantly cancel against other loops contributing to the one-loop $\gamma\gamma$ couplings of the A^0 and H^0, thereby leading to suppressed production rates. (Compare the rates of the high-tanβ scenarios, D_3^+ and SD_1^-, in Table III to those for lower tanβ scenarios with similar m_{A^0}.)

Of course, there are certainly SUSY scenarios that will yield viable A^0 and H^0 signals in the $b\bar{b}$, $t\bar{t}$, h^0Z, and h^0h^0 modes, in particular any model in which all SUSY states are more massive than one-half the Higgs boson masses. Nonetheless, we cannot ignore the fact that the very attractive dilatonlike boundary conditions suggested by superstring theory generally yield a sufficiently complex array of A^0 and H^0 decays as to make their detection in $\gamma\gamma$ collisions highly problematical.

We conclude that one should not count on being able to see the H^0 and A^0 in $\gamma\gamma$ collisions for integrated luminosities of order $L = 10\ \text{fb}^{-1}$ unless we become convinced by other experiments that the SUSY mass scale is quite high. This places increased onus on achieving much higher L or on building a machine with $\sqrt{s}$ sufficiently large that H^0A^0 and H^+H^- pair production will be possible via direct e^+e^- collisions. With regard to the latter, the gauge-coupling unified models typified by those explored here suggest that $\sqrt{s}$ above 500 GeV is generally required, with 1 TeV providing adequate energy for a large section of model parameter space. Of course, it remains to explore the degree to which SUSY decay modes and backgrounds complicate the detection of the above pair states [9].

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