Search for new Higgs physics at the Fermilab Tevatron

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(Received 15 November 2002; published 10 March 2003)

We discuss the Higgs boson mass sum rules in the minimal supersymmetric standard model (MSSM) in order to estimate the upper limits on the masses of top squarks as well as the lower bounds on the masses of the scalar Higgs boson state. The bounds on the scale of quark-lepton compositeness derived from the CDF Collaboration (Fermilab Tevatron) data and applied to the new extra gauge boson search are taken into account. These extra gauge bosons are considered in the framework of the extended $SU(2)_h \times SU(2)_l$ model. In addition, we discuss the physics of rare decays of MSSM Higgs bosons in both *CP*-even and *CP*-odd sectors and also some extra gauge bosons.

DOI: 10.1103/PhysRevD.67.055001 PACS number(s): 12.60.Jv, 14.70.Pw, 14.80.Cp, 14.80.Ly

I. INTRODUCTION

There are still some serious ingredients in the fundamental interactions of elementary particles that have not been experimentally verified. The Higgs particle(s) and new heavy neutral gauge bosons $G' \subset Z'$, $W^{\pm'}$, Z'', $W^{\pm''}$, etc., have not yet been established and the physics of those particles still remains elusive. In recent years, the amount of work searching for the Higgs and extra gauge bosons has increased considerably and this research subject is now one of the most exciting topics in searching for physics beyond the standard model (SM). On one hand, recent CERN e^+e^- collider LEP 2 experiments determined the lower bound of the Higgs boson mass to be approximately 114.1 GeV [1]. Furthermore, it should be pointed out that the Fermilab Tevatron data [2] for searching for the low energy effects of quark-lepton contact interactions on dilepton production taken at $\sqrt{s}=1.8$ TeV translate into lower bounds on the masses of extra neutral gauge bosons Z'. On the other hand, recent theoretical progress in the study of strong and electroweak interactions and the fundamental concept of natural extensions of the SM have changed the physics of interplay between the known matter fields and suggested exciting new phenomena related to new matter fields via new interactions.

One of the most challenging current topics beyond the SM is to study the physical implication of a set of Higgs particles and extra gauge bosons predicted by the minimal supersymmetric standard model (MSSM) in the Tevatron $\overline{p}p$ collider experiment at $\sqrt{s}=2$ TeV. The production of Higgs bosons or their decays (via heavy quark or lepton interactions) are promising and useful processes which could be searched for in the forthcoming high energy experiments. With the advent of run II at the Tevatron, it is expected to be shown that the intermediate heavy quark loop can be a sizable source of Higgs bosons in both *CP*-even (light h and heavy H) scalar and *CP*-odd pseudoscalar (A) sectors. As for extra neutral gauge bosons G', their existence is required by models addressing the physics beyond the SM, such as supersymmetry (SUSY), grand unification theory, superstrings and so on. The prediction for the masses $M_{G'}$ of the G' bosons is rather uncertain, although these bosons could be heavy enough with masses $M_{G'} \gg \mathcal{O}(m_Z)$ (m_Z is the mass of the Z boson). The simplest version of existing G' bosons is based on a model with extension of the standard SU(3) $\times$SU(2)$\times$U(1) gauge group by an extra U(1) factor [3]. There are also many other extended models of the SM, such as the ones in which the SU(2) gauge group is extended to SU(2)$\times$SU(2) [4–7] and to SU(2)$\times$SU(2)$\times$U(1) [8]. In those models, the massive SU(2) extra gauge bosons (corresponding to the broken generators) could couple to fermions in different generations with different strengths and thus could give the answer to the question of why the top quark is so heavy, since they might single out fermions of the third generation. Precision measurements of the electroweak parameters narrowed the allowed region of extra gauge boson masses, keeping Higgs boson masses consistent with radiative corrections including the supersymmetric ones.

One of the goals of the present work is to examine the potential of the run II experiment at the Tevatron in searching for h, H, and A Higgs bosons and new extra gauge bosons, and to give an estimation of the upper limits on the masses of top squarks as well as lower bounds on the masses of the scalar Higgs bosons in both the light and heavy mass sectors.

We first show how the existing Tevatron bounds on the scale of quark-lepton compositeness [9] can be adopted to provide an upper limit on the quantity $m_{\tilde{t}_1} \cdot m_{\tilde{t}_2}$, i.e., the product of the masses of top squark eigenstates $\tilde{t}_1$ and $\tilde{t}_2$. We also discuss how the lower bounds on the scalar Higgs boson masses can be obtained from the forthcoming Tevatron data. Furthermore, we emphasize that the forthcoming experiments for discovering supersymmetry in both the Higgs and quark sectors could lead to estimations of the masses of the neutral and charged extra gauge bosons Z' and

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$W^{\pm\prime}$, respectively. As an example, models in which the precision electroweak data allow the extra gauge bosons to be of order $\mathcal{O}(0.5\ \text{TeV})$ are, e.g., the noncommuting extended technicolor models [5]. Z' and $W^{\pm\prime}$ bosons with such masses are of interest, since they are within the kinematic reach of the Tevatron's run II experiments.

Then we also study the processes such as

$$\bar{p}p \to gg \to (h/H)X, AX, \tag{1}$$

keeping in mind that the main channels in the MSSM are the scalar ones, i.e., $\bar{l}l$, $\bar{Q}Q$, $\gamma\gamma h$, ggh (γ, g, l, and Q mean a photon, a gluon, a lepton, and a heavy quark, respectively) since the pseudoscalar mode is largely suppressed in the wide range of the MSSM parameters. Since in two-photon ($\gamma\gamma$) or two-gluon (gg) decays the invariant mass of the $\gamma\gamma$ or gg system would be identical to the mass of the decaying boson, a promising way to detect h and H bosons at the Tevatron is to search for the decays $h,H \to \gamma\gamma$ and $h,H \to gg$ as well. It is known that in high energy $\bar{p}p$ (or pp) collisions the contributions of gluonic interactions become large due to the increase of gluon densities in the proton. The Tevatron run II experiment could observe the lightest Higgs boson production via a two-gluon fusion $gg \to h$ with a cross section σ of order $\sigma \sim \mathcal{O}(1.0\ \text{pb})$ at $m_h \sim 110$ GeV [10]. Increasing the h boson mass up to 180 GeV would lead to decreasing of σ up to order $\mathcal{O}(0.2\ \text{pb})$. For A Higgs boson production the following important features are remarkable.

(i) The detection efficiency of the signal events has high accuracy because the decay of A Higgs bosons can be precisely modeled in the kinematic region for various decay channels, e.g., $A \to \bar{l}l, \gamma\gamma h, ggh$ (here, the leptons l run over electrons e, muons μ, and τ leptons).

(ii) The mass M_A of the A Higgs boson can be reconstructed from its final state to test the mass relation in the MSSM mass sum rule [11] at the tree level,

$$m_h^2 + M_H^2 = M_A^2 + m_Z^2, \tag{2}$$

with its deviation due to radiative corrections (m_h and M_H are the masses of the CP-even h and H Higgs bosons, respectively).

The outline of the article is as follows. In Sec. II, we define the model. Estimation of the upper limits of the masses of top squarks as well as the lower bounds on the masses of the scalar Higgs bosons will be discussed in Sec. III. Section IV is devoted to study of rare decays of h, H, and A Higgs bosons in the MSSM. Section V focuses on h Higgs boson production in the decay of an extra gauge boson Z_2. Finally, in Sec. VI, we give our conclusions.

II. THE EFFECTIVE MODEL

In the model of extended weak interactions governed by a pair of SU(2) gauge groups $SU(2)_h \times SU(2)_l$ for heavy (third generation) and light fermions (the labels h and l mean heavy and light, respectively) the gauge boson eigenstates are given by [12]

$$A^\mu = \sin\theta\ (\cos\phi\ W^\mu_{3_h} + \sin\phi\ W^\mu_{3_l}) + \cos\theta\ X^\mu \tag{3}$$

for a photon and

$$Z_1^\mu = \cos\theta\ (\cos\phi\ W^\mu_{3_h} + \sin\phi\ W^\mu_{3_l}) - \sin\theta\ X^\mu, \tag{4}$$

$$Z_2^\mu = -\sin\phi\ W^\mu_{3_h} + \cos\phi\ W^\mu_{3_l} \tag{5}$$

for the neutral gauge bosons Z_1, Z_2, respectively, which define the neutral mass eigenstates Z and Z' at the leading order of a free parameter x [9],

$$\begin{pmatrix} Z \\ Z' \end{pmatrix} \simeq \begin{pmatrix} 1 & \dfrac{-\cos^3\phi\ \sin\phi}{x\cos\theta} \\ \dfrac{\cos^3\phi\sin\phi}{x\cos\theta} & 1 \end{pmatrix} \cdot \begin{pmatrix} Z_1 \\ Z_2 \end{pmatrix}, \tag{6}$$

where θ is the usual weak mixing angle and ϕ is an additional mixing angle due to the existence of $SU(2)_h \times SU(2)_l$.

The parameter x in Eq. (6) is defined as the ratio $x = u^2/v^2$, where u is the energy scale at which the extended weak gauge group $SU(2)_h \times SU(2)_l$ is broken to its diagonal subgroup $SU(2)_L$, while $v \simeq 246$ GeV is the vacuum expectation value of the (composite) scalar field responsible for the symmetry breaking $SU(2)_L \times U(1)_Y \to U(1)_{e.m.}$ in the model of extended weak interactions. The generator of the $U(1)_{e.m.}$ group is the usual electric charge operator $Q = T_{3_h} + T_{3_l} + Y$. At large values of $\sin\phi$, the Z_2 boson has an enhanced coupling to the third generation fermions through the covariant derivative

$$\begin{aligned} D^\mu = \partial^\mu - i\ \frac{g}{\cos\theta}\ Z_1^\mu (T_{3_h} + T_{3_l} - \sin^2\theta \cdot Q) \\ -ig\ \ Z_2^\mu\ \left(-\frac{\sin\phi}{\cos\phi}\ T_{3_h} + \frac{\cos\phi}{\sin\phi}\ T_{3_l}\right). \end{aligned} \tag{7}$$

The Lagrangian density for an effective quark-lepton contact interaction looks like

$$\begin{aligned} \mathcal{L} \supset & \frac{1}{\Lambda_{LL}^2} [g_0^2 (\bar{E}_L \gamma_\mu E_L)(\bar{Q}_L \gamma^\mu Q_L) \\ & + g_1^2 (\bar{E}_L \gamma_\mu \tau_a E_L)(\bar{Q}_L \gamma^\mu \tau_a Q_L)] \\ & + \frac{g_e^2}{\Lambda_{LR}^2} (\bar{e}_R \gamma_\mu e_R)(\bar{Q}_L \gamma^\mu Q_L) \\ & + \left[\frac{1}{\Lambda_{LR}^2} (\bar{E}_L \gamma_\mu E_L) + \frac{1}{\Lambda_{RR}^2} (\bar{e}_R \gamma_\mu e_R) \right] \\ & \times \sum_{q:u,d} g_q^2 (\bar{q}_R \gamma^\mu q_R), \end{aligned} \tag{8}$$

where $E_L = (\nu_e, e)_L, Q_L = (u,d)_L$; g_i are the effective couplings and Λ_{ij} are the scales of new physics. The aim of the

Collider Detector at Fermilab (CDF) Collaboration analysis [2] was to search for deviation from the SM prediction in the dilepton production spectrum. If no such deviations are found, the lower bound of the Λ scale can be obtained. The embedding of the extra gauge bosons in the model beyond the SM gives rise to quark-lepton contact interactions in accordance with the following part of the Lagrangian density (see [9]):

$$\mathcal{L}\supset -\frac{g^2}{M_{Z'}^2}\left(\frac{\cot\phi}{2}\right)^2\left(\sum_{l:e,\mu}\overline{l_L}\gamma_\mu l_L\right)\left(\sum_{q:u,d,s,c}\overline{q_L}\gamma^\mu q_L\right), \tag{9}$$

where $g=e/\sin\theta$. We suppose that the couplings in the first two generations are the same in strength.

III. MASS BOUND ON TOP SQUARKS AND SOME HIGGS BOSON MASS ESTIMATIONS

In the MSSM, the mass sum rule (2) at the tree level is transformed into the following form because of the loop corrections [13]:

$$M_{Z'}=\frac{m_h^2-M_A^2+\delta_{ZZ'}-\Delta}{M_{Z'}+M_H}+M_H, \tag{10}$$

where $M_{Z'}$ is the mass of the Z' boson; $\delta_{ZZ'}=M_{Z'}^2-m_Z^2$. The correction Δ reflects the contribution from loop diagrams involving all the particles that couple with the Higgs bosons [14,15]

$$\Delta=\left(\frac{\sqrt{N_c}\ \ g m_t^2}{4\ \ \pi\ \ m_W\sin\beta}\right)^2\log\left(\frac{m_{\tilde{t}_1}\cdot m_{\tilde{t}_2}}{m_t^2}\right)^2, \tag{11}$$

where N_c is the number of colors, and m_t and m_W are the masses of the top quark and W boson, respectively. $\tan\beta$ defines the structure of the MSSM. The values of $\Delta\sim\mathcal{O}(0.01\ \mathrm{TeV}^2)$ have been calculated [14] for any choice of parameter space of the MSSM. We suggest that the measurement of $M_{Z'}$ would predict the masses of the mass eigenstates $\tilde{t}_1$ and $\tilde{t}_2$, since m_t and m_W are already measured in the experiments and m_h is restricted by the LEP 2 data [1] as $m_h<130$ GeV [16]; M_A and M_H are free parameters bounded by combined data coming from the MSSM parameter space and the experimental data [17].

Comparing Eqs. (8) and (9), one can get the relation between $M_{Z'}$ and Λ as [9]

$$M_{Z'}=\sqrt{\alpha_{e.m.}}\ \ \Lambda\cot\phi/(2\sin\theta), \tag{12}$$

where the value of Λ was constrained from the CDF data at $\sqrt{s}=1.8$ TeV as $\Lambda>3.7$ TeV or 4.1 TeV, depending on the contact interactions for the left-handed electrons or muons, respectively, found at the 95% confidence level [2,9].

In the decoupling regime of the MSSM Higgs sector where the couplings of the light *CP*-even Higgs boson h in the MSSM are identical to those of the SM Higgs bosons and thus the *CP*-even mixing angle α behaves as $\tan\alpha\to-\cot\beta$ with the $M_A\gg m_Z$ relation, one can get $M_H^2\simeq M_A^2+m_Z^2\sin^2(2\beta)+\mu^2$, which leads to disappearance of the H Higgs boson mass in Eq. (10). Here, μ is the positive massive parameter which can, in principle, be defined from an experiment searching for separation of two degenerate heavy Higgs bosons A and H. This behavior, verified at the tree level, continues to hold even when radiative corrections are included. It has been checked that this decoupling regime is an effective one for all values of $\tan\beta$ and that the pattern of most of the Higgs couplings results from this limit.

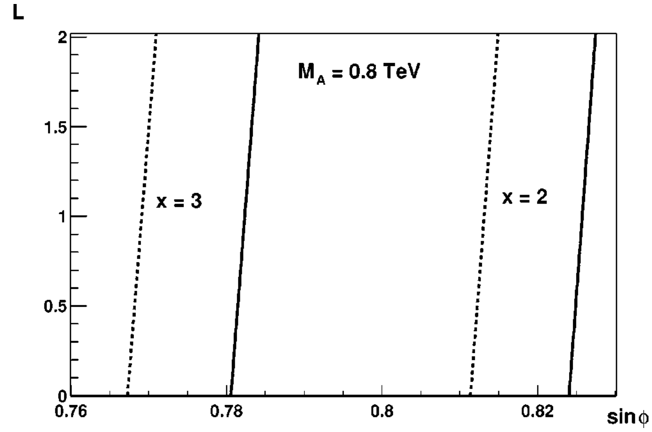

FIG. 1. The upper limit on $L\equiv\log[(m_{\tilde{t}_1}\cdot m_{\tilde{t}_2})/m_t^2]$ as a function of $\sin\phi$ for different values of $x=2$ and 3; $\mu=m_h=120$ GeV (dashed line), $\mu=0$ (solid line) for $M_A=0.8$ TeV; $\tan\beta=30$.

In studying the mass relation (10) from the extended electroweak gauge structure, we must be aware of the issues related to the structure of $M_{Z'}$ in both sides of Eq. (10). We suppose that $M_{Z'}$ in the left-hand side (LHS) of Eq. (10) is the mass to be determined using the CDF analysis data [2,9]. Therefore, one can approximate its mass via the phenomenological relation (12) while the RHS of Eq. (10) is model dependent, where, to leading order, the mass $M_{Z'}$ in the extended weak interaction model is $M_{Z'}=m_W\ \sqrt{x}/(\cos\phi\sin\phi)$ [9] in the region where $\cos\phi<\sin\phi$. With the help of the CDF restriction for Λ [2] entering into Eq. (12), one can easily find the upper limit of the product of $m_{\tilde{t}_1}\cdot m_{\tilde{t}_2}$ from the following relation [13]:

$$\Delta<(B+M_H^\star)(B-fC)+m_h^2-m_Z^2\ \ (1-\sin^2 2\beta)+\mu^2, \tag{13}$$

where $M_H^\star=(M_A^2+m_Z^2\sin^2 2\beta+\mu^2)^{1/2}$, $f\equiv f(\phi)=\cot\phi\ \sqrt{\alpha_{e.m.}}/(2\sin\theta)$, $B\equiv B(x,\phi)=m_W\sqrt{x}/(\cos\phi\sin\phi)$, and C is the minimal value of the Λ scale extracted from the CDF analysis [2]. The masses of Z and W bosons are currently known with errors of a few MeV each [18], whereas the mass of the top quark is known with errors of a few GeV [18]. Note that the dependence of particle couplings via $\tan\beta$ enters into the radiative correction Δ in Eq. (11) and the mass M_H defined in the decoupling regime. Thus, the upper limit on $m_{\tilde{t}_1}\cdot m_{\tilde{t}_2}$ can be accurately predicted by precision measurements of the lower bound of $M_{Z'}$ and the masses of the Higgs bosons h and A.

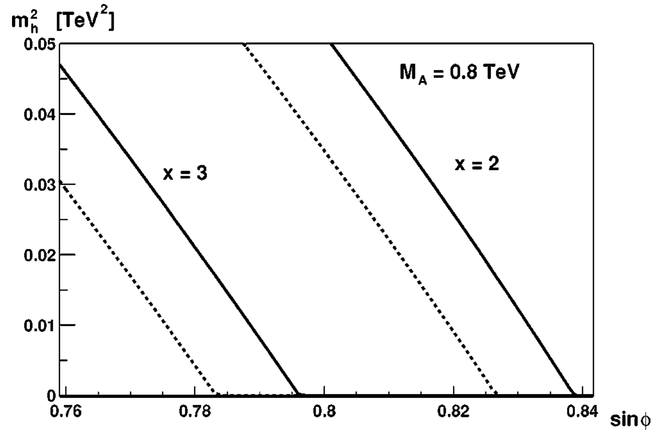

FIG. 2. The lower bound on m_h^2 as a function of $\sin\phi$ for different values of $x=2$ and 3; $\mu=m_h=120$ GeV (dashed line), $\mu=0$ (solid line) for $M_A=0.8$ TeV; $\tan\beta=30$. The regions of the parameter space lying above a given line are allowed by the present model.

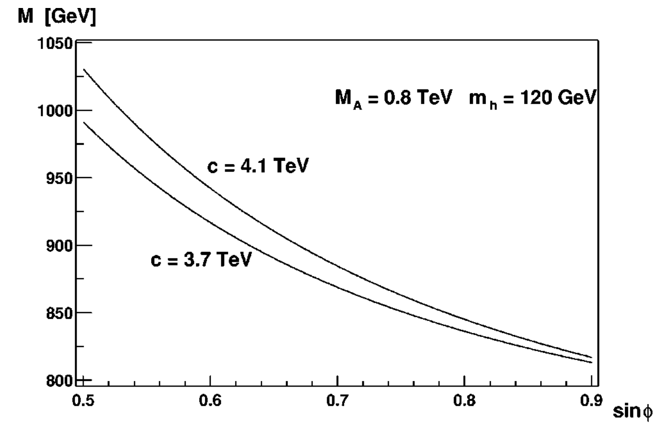

FIG. 3. The lower bound on $M\equiv(\Sigma_{j=1}^2 M_{H_j}^2)^{1/2}$ as a function of $\sin\phi$.

Figure 1 shows the upper limit on $L\equiv\log[(m_{\tilde{t}_1}\cdot m_{\tilde{t}_2})/m_t^2]$ as a function of $\sin\phi$ for $x=2$ and $x=3$ at fixed values of μ and M_A. We use the range of the mixing parameter $0.75\leqslant\sin\phi\leqslant0.85$ for a Z' boson [9] where the luminosity required to exclude SU(2) Z' bosons of various masses is lowest. The corresponding range of the lower bound on x is $3.9\geqslant x\geqslant1.6$ (left-handed muons and up-type quarks are taken into account) for the $\sin\phi$ range above mentioned. The regions of the parameter space lying below a given line are allowed by the present model. At present, the LEP bounds on the mass of the A Higgs boson are $M_A>88.4$ GeV [1]. This result corresponds to the large $\tan\beta$ region. We see that the function L is rather sensitive within the changing of $\sin\phi$, i.e., the ratio of gauge couplings g/g_l. Here, $g^{-2}=g_l^{-2}+g_h^{-2}$, where g_l is associated with the $SU(2)_l$ group and defines the couplings to the first and second generation fermions, whose charges under subgroup $SU(2)_l$ are the same as in the SM, while g_h originates from the $SU(2)_h$ group, which governs the weak interactions for the third generation (heavy) fermions. In the range of $\sin\phi$ presented in Fig. 1, the width $\Gamma_{Z'}$ of the Z' boson falls to a minimum in the neighborhood of $\sin\phi=0.8$ [9], due to the decreasing couplings to the first two generations of fermions. In the range $\sin\phi>0.8$, $\Gamma_{Z'}$ grows large, due to the rapid growth in the third generation coupling.

The CDF analysis of the contact interaction between left-handed muons and up-type quarks is taken into account ($C=4.1$ TeV) in our calculations. The lower bounds of m_h^2 are illustrated in Fig. 2. The constraints are given for different ratios of $x=2$ and 3 as a function of $\sin\phi$. In our calculations, the parameters of the model are typically chosen as $m_h=120$ GeV, $\tan\beta=30$. The value of the neutral *CP*-odd Higgs boson A mass is set to be $M_A=0.8$ TeV within the typical upper limit of M_A kinematically allowed at the Tevatron run II energy. Herewith, we suppose that A Higgs boson can be identified, e.g., via two-lepton decays in the AZ associated production process $p\bar{p}\to AZ+X$ with $\sqrt{s}>M_A+m_Z$. Here, we did not use the mass difference between $\tilde{t}_1$ and $\tilde{t}_2$ mass eigenstates, and we set $m_{\tilde{t}_1}=m_{\tilde{t}_2}=1$ TeV (see Fig. 2). The regions of the parameter space lying above a given line are allowed by the present data.

In more extended SUSY models, their mass sum rules can give some useful estimations with the help of the CDF data [2]. For example, in the minimal E_6 superstring theory, the particle spectrum consists of three scalar Higgs bosons h,H_1,H_2, a pseudoscalar Higgs A, a charged Higgs boson pair $H^{\pm}$, and two neutral gauge bosons Z and Z'. Among these particles, there exists a mass sum rule, at the tree level, of the form [19]

$$M_{Z'}^2=m_h^2+M_{H_1}^2+M_{H_2}^2-M_A^2-m_Z^2. \tag{14}$$

The analytical expressions for the loop corrections are unknown yet. However, it is known that the one-loop corrections can be summarized into a term logarithmically dependent on the SUSY sector mass scale [19]. Considering that m_h can be identified with the lower bound on the Higgs boson mass, we obtain the lower bound on the sum $M_{H_1}^2+M_{H_2}^2$ at fixed M_A as a function of $\sin\phi$:

$$\sum_{j=1}^{2} M_{H_j}^2>M_A^2+m_Z^2-m_h^2+\frac{\alpha_{e.m.}\,C^2\cot^2\phi}{4\sin^2\theta}. \tag{15}$$

The results of the calculation of the lower bound on $M\equiv(\Sigma_{j=1}^2 M_{H_j}^2)^{1/2}$ as a function of $\sin\phi$ are given in Fig. 3.

We have used the scales of new physics $\Lambda>C$ coming from the CDF analysis [9] at 95% confidence level in which contact interactions were assumed only between left-handed electrons (muons) and up-type quarks: $\Lambda>4.1$ TeV and $\Lambda>3.7$ TeV for left-handed muons and left-handed electrons, respectively.

IV. RARE DECAYS OF THE HIGGS BOSONS

As is well known, Higgs bosons dominantly couple to heavy particles even in the MSSM. Observation of h, H, and A Higgs bosons depends on the model parameters including the masses m_h, m_H, and M_A.

A. The decays $h,H \to gg,\gamma\gamma$

Let us begin with study of the decays $X \to gg$ and $X \to \gamma\gamma$ $(X=h,H)$, where the former dominates over the latter. The decay width of $X \to gg$ (the radiative corrections are not included) is given by

$$\Gamma(X\to gg)=\frac{g^2\ \alpha_s^2\ m_X^3}{128\ \pi^3\ m_W^2}\cdot|\tilde{F}_Q|^2, \tag{16}$$

where the transition form factor $\tilde{F}_Q$ is provided by an intermediate quark loop containing b and t quarks [20,21]:

$$\begin{aligned}\tilde{F}_Q=\rho_{Xb}\ \ \tau_b\left\{\frac{(\tau_b-1)}{4}\left[\pi^2-\ln^2\left(\frac{1+\xi_b}{1-\xi_b}\right)\right]-1\right\}\\ +\rho_{Xt}\ \ \tau_t\left\{(\tau_t-1)\left[\arcsin^2\left(\frac{m_X}{2\ \ m_t}\right)\cdot\theta(2\ \ m_t-m_X)\right.\right.\\ \left.\left.+\frac{1}{4}\ \left[\pi^2-\ln^2\left(\frac{1+\xi_t}{1-\xi_t}\right)\right]\cdot\theta(m_X-2m_t)\right]-1\right\}.\end{aligned} \tag{17}$$

Here, $\tau_Q=(2\ m_Q/m_X)^2$, $\xi_Q=\sqrt{1-\tau_Q}$, $\rho_{hQ}=-\sin\alpha/\cos\beta$, and $\rho_{HQ}=\cos\alpha/\cos\beta$ for b quarks; $\rho_{hQ}=\cos\alpha/\sin\beta$ and $\rho_{HQ}=\sin\alpha/\sin\beta$ for t quarks; the mixing angle α diagonalizes the CP-even Higgs squared-mass matrix. For comparison, the decay width $\Gamma(X\to\gamma\gamma)$ can easily be obtained from the corresponding decay to two gluons (16) by a simple replacement $\alpha_s \to (3/\sqrt{2})\alpha\ e_Q^2$ (the heavy quark charges e_Q are included in the form factors) and $\tilde{F}_Q \to F_Q+F_W+F_H^{\pm}$:

$$\Gamma(X\to\gamma\gamma)=\frac{g^2\ \alpha^2\ m_X^3}{1024\ \pi^3\ m_W^2}\cdot|F_Q+F_W+F_{H^\pm}|^2, \tag{18}$$

where the form factor F_Q contributed from heavy quarks Q is given by

$$\begin{aligned}F_Q=2\left\{\rho_{Xb}\frac{\tau_b}{3}\left[\frac{(\tau_b-1)}{4}\left(\pi^2-\ln^2\frac{1+\xi_b}{1-\xi_b}\right)-1\right]\right.\\ +\frac{4}{3}\ \rho_{Xt}\ \ \tau_t\left[(\tau_t-1)\left(\arcsin^2\left(\frac{m_X}{2\ \ m_t}\right)\cdot\theta(2\ \ m_t-m_X)\right.\right.\\ \left.\left.\left.+\frac{1}{4}\ \left(\pi^2-\ln^2\frac{1+\xi_t}{1-\xi_t}\right)\cdot\theta(m_X-2\ \ m_t)\right)\right]-1\right\},\end{aligned} \tag{19}$$

and the F_W from weak bosons $W^\pm$ and $F_{H^\pm}$ from charged Higgs bosons $H^\pm$ are [22]

$$\begin{aligned}F_W=\rho_{XW}\left\{2+3\tau_W+3\tau_W(2-\tau_W)\right.\\ \times\left[\arcsin^2\left(\frac{m_X}{2\ \ m_W}\right)\cdot\theta(2\ \ m_W-m_X)\right.\\ \left.\left.+\frac{1}{4}\ \left(\pi^2-\ln^2\frac{1+\xi_W}{1-\xi_W}\right)\cdot\theta(m_X-2\ \ m_W)\right]\right\}\end{aligned} \tag{20}$$

and

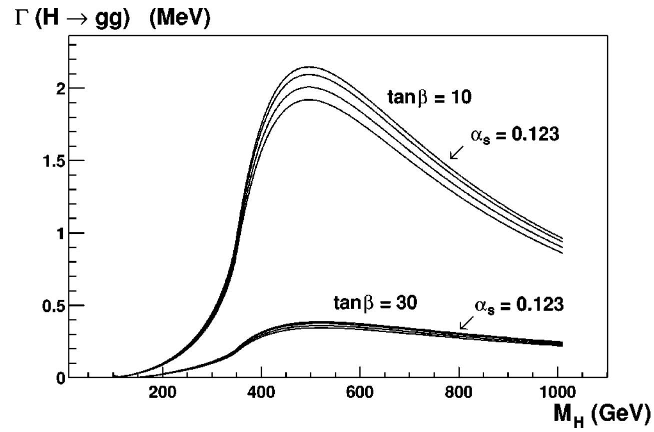

FIG. 4. The decay width $\Gamma(H\to gg)$ as a function of M_H for different α_s (from the bottom $\alpha_s=0.110$, 0.115, 0.119, 0.123) and $\tan\beta=10$ and 30.

$$F_{H^\pm}=\rho_{XH^\pm}\left(\frac{m_W}{M_{H^\pm}}\right)^2\left(1-\tau_{H^\pm}\arcsin^2\frac{m_X}{2\ \ M_{H^\pm}}\right), \tag{21}$$

respectively. Here, $\rho_{hW}=\sin(\beta-\alpha)$, $\rho_{HW}=\cos(\beta-\alpha)$, $\rho_{hH^\pm}=\rho_{hW}+[\cos 2\beta\sin(\beta+\alpha)]/2\cos^2\theta_W$, and $\tau_{H^\pm}=(2\ M_{H^\pm}/m_X)^2$. The calculated results for $\Gamma(H\to gg)$,$\Gamma(h\to gg)$ and $\Gamma(H\to\gamma\gamma)$,$\Gamma(h\to\gamma\gamma)$ are given in Figs. 4, 5 and 6, 7, respectively, where the mass of the CP-odd A Higgs boson was typically taken to be $M_A=0.8$ TeV.

In our calculations, the couplings between h Higgs bosons and left- and right-handed fermions as well as between h Higgs bosons and charginos were neglected. The following relation between the mixing angle α and $\tan\beta$:

$$\cot\alpha=-\tan\beta\left[1+2\left(\frac{m_Z}{M_A}\right)^2\cos 2\beta\right]+\mathcal{O}\left(\frac{m_Z^4}{M_A^4}\right) \tag{22}$$

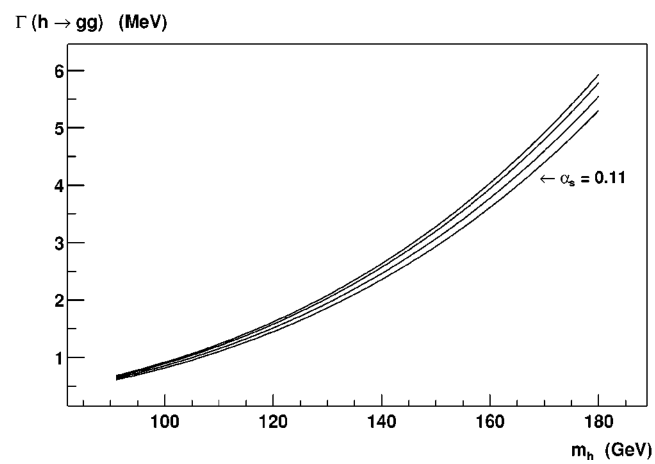

FIG. 5. The decay width $\Gamma(h\to gg)$ as a function of m_h for different α_s (from the bottom $\alpha_s=0.110$, 0.115, 0.119, 0.123) in the decoupling limit.

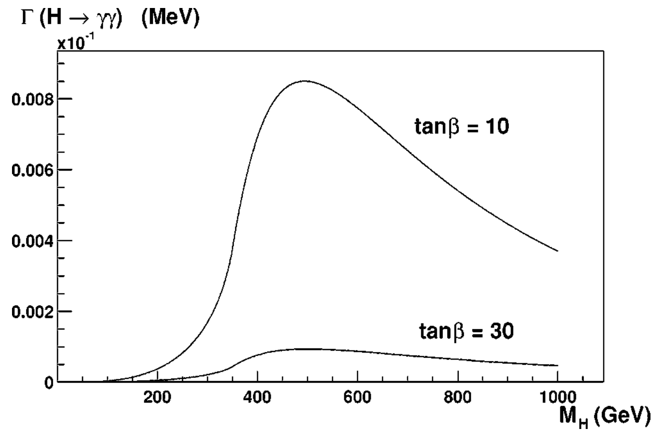

FIG. 6. The decay width $\Gamma(H\to\gamma\gamma)$ as a function of M_H for different $\tan\beta=10$ and 30.

is useful in evaluating h Higgs boson couplings to b and t quarks in the decoupling regime.

Here, we point out that the run II at the Tevatron could observe the Higgs boson h through the promising final states of the gg and $\gamma\gamma$ channels or the $\bar{b}b$ final state.

B. The decays $A\to\bar{l}l$

In the MSSM, the decay amplitude of a *CP*-odd neutral Higgs boson A (with mass M_A and four-momentum P_μ) into a lepton l and antilepton $\bar{l}$ with four-momenta p_μ and $(P-p)_\mu$, respectively, is given by the following Feynman amplitude:

$$\mathrm{Am}(A\to\bar{l}l)=i\sum_{Q:b,t}\bar{u}(p)\;\gamma_5\;F_Q(M_A^2)v(P-p),\tag{23}$$

where $F_Q(t=M_A^2)$ is the complex function (form factor) providing the transition from the A Higgs boson into a lepton pair. One can consider this transition via the intermediate off-shell spin-1 bosons V^* (photons, $W^\pm$ bosons, and so on) within the vertex AV^*V^*:

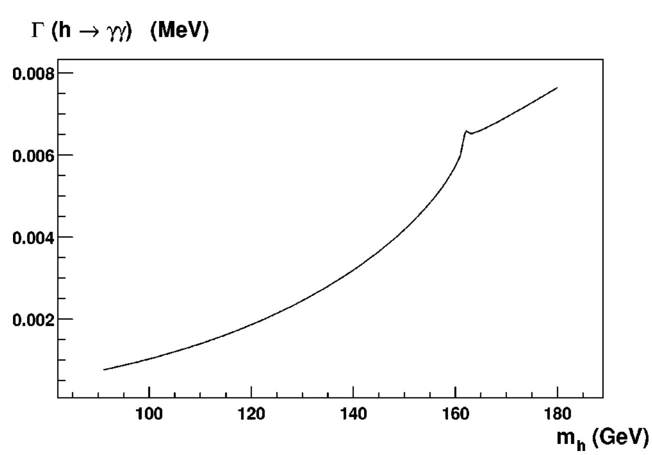

FIG. 7. The decay width $\Gamma(h\to\gamma\gamma)$ as a function of m_h in the decoupling limit.

$$\Gamma_{\mu\nu}=\epsilon_{\mu\nu\alpha\beta}\,k^\alpha\,(P-k)^\beta\,F_A[k^2,(P-k)^2]\tag{24}$$

with

$$F_A[k^2,(P-k)^2]=\sum_{Q:b,t}\rho_Q\;e_Q^2\;f_{VV}^Q\;\hat{F}_{QA}[k^2,(P-k)^2].\tag{25}$$

Here, $\rho_Q=\tan\beta$ for the b-quark loop and $\rho_Q=\cot\beta$ for the t-quark loop; k_μ is the four-momentum of a V^* boson. There is a natural normalization condition $F_A(0,0)=f_{VV}$ in Eq. (24) where the decay constant f_{VV} is given for the decay of an A Higgs boson into two real V bosons. The form of F_A in Eq. (25) should be taken so that the Feynman integrals are not divergent. Here, we simply take $\hat{F}_{QA}[k^2,(P-k)^2]$ as

$$\hat{F}_{QA}[k^2,(P-k)^2]=\frac{\lambda_Q^2}{\lambda_Q^2-k^2-(P-k)^2}\tag{26}$$

which has good analytical properties, and carries both the vector dominance features and the properties of the right static limit. The parameter λ_Q must be large enough, and for numerical estimations we put $\lambda_Q^2=M_A^2$ where $M_A\sim\mathcal{O}(1\text{ TeV})$. To be more precise in the calculations of the decay width of the process $A\to\bar{l}l$, one has to take into account the fact that the form factor $F_A[k^2,(P-k)^2]$ should be a complex function with the following real part:

$$\mathrm{Re}\,F_A(M_A^2)=\frac{P}{\pi}\int_0^\infty\frac{\mathrm{Abs}[F_A(t)]}{t-M_A^2}\;dt.\tag{27}$$

Here, the absorptive part $\mathrm{Abs}[F_A(t)]$ can be found if we take the S_0^1 leptonic pair in the final state and take into account the unitarity condition for the decay amplitude:

$$\mathrm{Abs}[F_A(t)]=\frac{1}{2\;i\;\sqrt{2\;t}}\int d\Omega_I\;\delta^4\left(P-\sum_I q_I\right)\times\langle I|T|(\bar{l}l)_{S_0^1}\rangle^*\;\mathrm{Am}(A\to I).\tag{28}$$

In Eq. (28), the integration over the phase-space volume Ω_I for all possible intermediate states I, i.e., the off-shell $\gamma^*\gamma^*$ state, the heavy quark–antiquark state, the heavy quark, and the off-shell γ^* quantum are taken into account. The decay width $\Gamma(A\to\bar{l}l)$ normalized to two-photon decays is calculated as follows [20]:

$$\mathrm{Br}(A\to\bar{l}l/\gamma\gamma)\equiv\frac{\Gamma(A\to\bar{l}l)}{\Gamma(A\to\gamma\gamma)}=2\;\xi_l\;\frac{1}{(4\;\pi)^2}\left(\frac{m_l}{M_A}\right)^2\frac{1}{\pi^2}\;|R|^2\tag{29}$$

with

TABLE I. The values $\mathrm{Br}(A\to \bar{l}l/\gamma\gamma)$ as a function of M_A where $l=\mu^-$ and τ^-.

	M_A (TeV)					
	0.10	0.25	0.50	0.75	1.00	1.20
$\mathrm{Br}(A\to\mu^+\mu^-/\gamma\gamma)\times 10^8$	3.00	1.00	0.28	0.15	0.10	0.007
$\mathrm{Br}(A\to\tau^+\tau^-/\gamma\gamma)\times 10^6$	0.84	0.27	0.11	0.06	0.04	0.02

$$R=\frac{1}{\xi_l}\left[\frac{i\ \pi}{2}\ln\frac{1-\xi_l}{1+\xi_l}+\frac{1}{4}\ln^2\frac{1+\xi_l}{1-\xi_l}-\ln\frac{1+\xi_l}{1-\xi_l}+\frac{\pi^2}{12}-\Phi\left(\frac{1-\xi_l}{1+\xi_l}\right)\right], \tag{30}$$

where $\Phi(x)=\int_0^x dt\ \ln(1+t)/t$, $\xi_l=\sqrt{1-(2\ m_l/M_A)^2}$. The normalized decay rate (29) is very convenient because the dependence on $\tan\beta$ is canceled for the VV channel. The corresponding SM background comes from the Drell-Yan production

$$\bar{q}q\to\gamma^*,Z^*,Z'^*\to\bar{l}l.$$

The signals we are looking for would be identified by the peaks of the invariant mass of $\tau^+\tau^-$ and/or $\mu^+\mu^-$ pairs. In Table I we show the calculated results of the decay widths $A\to\mu^+\mu^-$ and $A\to\tau^+\tau^-$ normalized to $\gamma\gamma$- channel as the function of M_A. At the end of this part, we conclude that, apart from reducing the $A\to\tau^+\tau^-$ signal due to some experimental constraints, this decay mode could be detected at the Tevatron's run II (almost) as easily as the corresponding signal of $A\to\mu^+\mu^-$ decay.

C. The decays $A\to\gamma\gamma h, ggh$

In addition to the proposal for searching for A Higgs bosons via the decay $A\to Zh$ [10], we suggest investigating the decays $A\to\gamma\gamma h$ and $A\to ggh$ which can be relevant at the Tevatron's run II for $M_A>m_h$ at large $\tan\beta$. We first note that the rate of the $\gamma\gamma h$ channel might be more promising for discovery of A Higgs bosons than those of the $\tau^+\tau^-$ and $\mu^+\mu^-$ channels.

TABLE II. $\mathrm{Br}(A\to\gamma\gamma h/\gamma\gamma)\times 10^{-2}$ as a function of κ for $\tan\beta=5,10,20$, and 50; M_A is taken to be $M_A=0.8$ TeV as a typical value.

	$\kappa=m_h/M_A$					
	0.05	0.1	0.2	0.3	0.4	0.5
$\tan\beta=5$	0.066	0.054	0.044	0.032	0.025	0.014
$\tan\beta=10$	0.26	0.22	0.18	0.13	0.10	0.058
$\tan\beta=20$	1.09	0.91	0.73	0.54	0.42	0.24
$\tan\beta=50$	6.54	5.43	4.43	3.28	2.54	1.43

The matrix element of the decay of A Higgs bosons with momentum P into γ, γ, and h with momenta k_1,k_2, and k_3, respectively, has the following form:

$$M(A\to\gamma\gamma h)=\frac{\sqrt{N_c}\ \pi\ \alpha\ ig}{m_W\ \sqrt{M_A}}\times\mathrm{Tr}\left[\sum_{Q:b,t}\gamma_5\ \rho_Q\ e_Q^2\ m_Q\times(M_A-\hat{P})\Gamma_Q\ \rho_{hQ}\ T_Q\right], \tag{31}$$

where T_Q is the amplitude providing the transition of the $\bar{Q}Q$ virtual quark pair into the $\gamma\gamma h$ final state (the permutations are taken into account); Γ_Q is the vertex function describing the couplings $h\bar{Q}Q$ taking into account the ρ_Q flavor-dependent factor. The differential distribution of the decay width $A\to\gamma\gamma h$ over the invariant mass $\tilde{s}$ of the final state's two-photon pairs and normalized to the $\gamma\gamma$ width is given by

$$\frac{1}{\Gamma(A\to\gamma\gamma)}\frac{d\Gamma(A\to\gamma\gamma h)}{d\tilde{s}}=\rho_{hQ}^2\frac{G_F\ a^3\ \tilde{s}\ M_A^2}{8\ \sqrt{2}\ \pi^2\ (a^2-\tilde{s}^2)^2\ (a+\tilde{s}/2)}, \tag{32}$$

where $a=(1-\kappa^2)/2$, $\kappa=m_h/M_A$, $\tilde{s}=s/M_A^2$, $s=(k_1+k_2)^2$. In Table II, we present $\mathrm{Br}(A\to\gamma\gamma h/\gamma\gamma)$ as a function of κ for different values of $\tan\beta=5,10,20$, and 50 at fixed $M_A=0.8$ TeV.

The width $\Gamma(A\to ggh)$ can be obtained from the corresponding decay width for $A\to\gamma\gamma h$ by making the replacement $\alpha\ e_Q^2\to(\sqrt{2}/3)\ \alpha_s$. The relative decay width of $A\to ggh$ normalized to $A\to gg$ gives an identical value to the one listed in Table II at $M_A=0.8$ TeV.

Compared with the $A\to\tau^+\tau^-/\mu^+\mu^-$ cases as discussed in Sec. IV B, we conclude that the $A\to\gamma\gamma h$ channel in the MSSM could be observed for $0.15\leqslant\kappa\leqslant 0.2$ (120 GeV$\leqslant m_h\leqslant$160 GeV) and for $\tan\beta>5$.

V. THE DECAYS $Z_2\to Z_1H$

To reach a unified theory of all interactions one could start with the grand unification theory (GUT) group SU(5) [23] which is the minimal unification group of strong and electromagnetic interactions. However, this example was ruled out by several experiments such as searching for the proton decay. The natural extension leads to SO(10) [24]. It is known that all unification groups larger than SU(5) have extra gauge bosons, the neutral Z' bosons and the charged $W^{\pm\prime}$ ones. Experimental signals of those extra gauge bosons would give very important information about the underlying GUT and its origin. The search for Z' and $W^{\pm\prime}$ is therefore an important subject for the physical program at the Tevatron's run II, where very high precision measurements can give valuable information on Z' and $W^{\pm\prime}$ because they are sensitive to rare processes, including their decays.

The cross section $\sigma(\bar{p}p\to Z'\bar{l}l)$ for Z' production at the Tevatron is inversely proportional to the total decay width $\Gamma_{Z'}$. Obviously, the inclusion of extra channels in Z' decays leads to $\Gamma_{Z'}$ becoming larger and $\sigma(\bar{p}p\to Z'\bar{l}l)$ smaller. As already pointed out [25,26], if the decays of Z' into $h+\{\bar{Q}Q\}_{s=1},h\bar{l}l$ are kinematically allowed, searches for h Higgs bosons or exotic heavy quark–antiquark bound states $\{\bar{Q}Q\}_{s=1}$ with spin $s=1$ would be interesting channels. The study of the decay $Z'\to h+\{\bar{Q}Q\}_{s=1}$ [25,26] would provide useful information about the nature of the extended gauge structure such as the couplings of Z' with heavy quarks in both the vector and axial-vector sectors, the couplings of h with heavy quarks, $Z-Z'$ mixing effects, the physics of mass eigenstates Z_1 and Z_2, and interplay with the matter fields.

In an effective rank-5 model, including only one extra neutral gauge boson Z', the interaction Lagrangian is given in the standard manner:

$$-\mathcal{L}_{int}=\frac{1}{2}g_1 Z_\mu\left[\sum_f \bar{\Psi}_f\gamma^\mu(g_v^f-g_a^f\gamma_5)\Psi_f\right] + \frac{1}{2}g_2 Z'_\mu\left[\sum_f \bar{\Psi}_f\gamma^\mu(g_v^{f\prime}-g_a^{f\prime}\gamma_5)\Psi_f\right], \quad (33)$$

where Ψ_f is the fermion field with flavor *f*. The first term in Eq. (33) is written down within the SM where $g_v^f=T_{3\,L}-2\,Q_f\sin^2\theta_W$, $g_a^f=T_{3\,L}$; $g_1=g/\cos\theta_W$ is the SM coupling constant, and $T_{3\,L}$ and Q_f are the third component of the weak isospin and the electric charge, respectively. The pairs (g_v^f,g_a^f) and $(g_v^{f\prime},g_a^{f\prime})$ represent the chiral properties of interactions of Z and Z' bosons with Ψ_f, respectively. The mass eigenstates Z_1 and Z_2 are parametrized by a mixing angle θ originating from the mixing between weak eigenstates Z and Z':

$$\begin{pmatrix} Z_1 \\ Z_2 \end{pmatrix}=\begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix}\cdot\begin{pmatrix} Z \\ Z' \end{pmatrix}. \quad (34)$$

Therefore, the Lagrangian (33) is replaced by

$$\mathcal{L}_{int}=\frac{-g}{2\cos\theta_W}\sum_f [Z_{1_\mu}\bar{\Psi}_f\gamma^\mu(V^f-A^f\gamma_5)\Psi_f + Z_{2_\mu}\bar{\Psi}_f\gamma^\mu(V^{f\,\prime}-A^{f\,\prime}\gamma_5)\Psi_f], \quad (35)$$

where

$$V^f=g_v^f\cos\theta+\frac{g_2}{g_1}\ g_v^{f\prime}\sin\theta,\quad A^f=g_a^f\cos\theta+\frac{g_2}{g_1}\ g_a^{f\prime}\sin\theta \quad (36)$$

and

$$V^{f\prime}=\frac{g_2}{g_1}\ g_v^{f\prime}\cos\theta-g_v^f\sin\theta,\quad A^{f\prime}=\frac{g_2}{g_1}\ g_a^{f\prime}\cos\theta-g_a^f\sin\theta \quad (37)$$

with $g_2=g_1\sqrt{(5/3)\ \sin^2\theta_W\,\lambda}\simeq g_1\cdot 0.62\sqrt{\lambda}$, $\lambda\sim\mathcal{O}(1)$ [27]. We use the LEP measured value $\sin^2\theta_W(\overline{MS})=0.23117$ [18].

The decays of Z_2 are a promising place to search for a *CP*-even light Higgs boson h. Here, the effects of heavy quarks cannot be neglected. As a result, there exists an effective h-gluon-gluon interaction which arises from the triangle diagram with a heavy quark loop and does not decouple in the limit of large quark masses. One can separate out the heavy quark contribution and use an effective low energy theorem [28–30] for the Higgs boson interactions. In the limit $M_{Z_2}\gg m_h$ [$m_h\sim\mathcal{O}(100\text{ GeV})$, M_{Z_2} is the mass of Z_2] when h is a constant field, the interaction of h is reproduced by rescaling all the mass terms $m_j=m_j(1+h/v)$, j = quarks, $W,Z,Z',\ldots,$ and $\alpha_s\to\alpha_s+\delta\alpha_s$ with $\delta\alpha_s=\alpha_s^2\ h/(3\ \pi\ v)$ [28–30]. Here, the number of heavy quarks is restricted to one, that is, only the top quark loop will be involved in the game. The interaction Lagrangian becomes

$$\mathcal{L}_{int}=(1+h/v)\rho_{ht}m_t\bar{t}t+\frac{\alpha_s\ \rho_{ht}}{12\pi v}\ hG_{\mu\nu}G^{\mu\nu} -(1+h/v)^2\left(m_W^2W_\mu^+W^\mu+\frac{1}{2}m_Z^2Z_\mu Z^\mu + \frac{1}{2}M_{Z'}^2Z'_\mu Z^{\mu\,\prime}\right), \quad (38)$$

where $G_{\mu\nu}$ is the standard gluon field strength tensor and *hgg* interactions are induced by the top quark loop.

If the mass of the h Higgs boson is low enough for the decay $Z_2\to Z_1 h\to\bar{l}lh$ to be kinematically allowed, the Tevatron bounds on this transition could severely constrain the structure of h couplings. The relative differential distribution of the decay width $\Gamma(Z_2\to\bar{l}lh)$ over the dimensionless variable $x=(p_l+p_{\bar{l}})^2/M_{Z_2}^2$ (p_l and $p_{\bar{l}}$ are the momenta of a lepton and antilepton, respectively) is given by

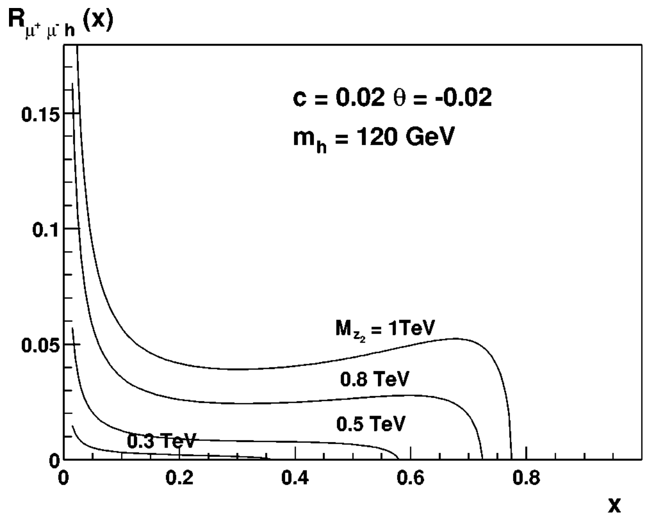

FIG. 8. The $x=(p_l+p_{\bar{l}})^2/M_{Z_2}^2$ distribution of $R_{\mu^+\mu^-h}$ in the decay $Z_2\to\mu^+\mu^-h$ for $M_{Z_2}=0.3$ TeV, 0.5 TeV, 0.8 TeV, and 1 TeV.

$$R_{\bar{l}lh}(x)\equiv\frac{1}{\Gamma(Z_2\to\bar{l}l)}\frac{d\ \ \Gamma(Z_2\to\bar{l}lh)}{dx}$$

$$=\frac{g^2\ \ V'VM_{Z_2}^2\ \ \rho_{ht}^2}{4\cos^2\theta_W\ \ \overline{V'}\ \ 24\ \ \pi^2\ \ v^2\ \ x}[(1-a_h)^2+x^2$$

$$-2x(1+a_h)]^{1/2}\left(1-\frac{4}{x}a_l\right)^{1/2}$$

$$\times\left(1+\frac{2}{x}a_l\right)\frac{(1-a_h)^2+x^2+2x(2-a_h)}{(1-x)^2+c^2},\quad(39)$$

where $V'=0.62\ \lambda^{1/2}-\theta\ g_v$, $V=g_v+0.62\ \lambda^{1/2}\ \theta$ [27], $\overline{V'}=(\sqrt{3}/2)(1-4\sin^2\theta_W)^{1/2}$ [31], $a_j=(m_j/M_{Z_2})^2$ with $j=l,h$, and $c=\Gamma_{Z_2}/M_{Z_2}$. To get $\overline{V'}$ we have taken into account the same normalization for the Z_2-lepton couplings as the usual one for the Z-lepton interplay within the SM. In the formula (39), the Drell-Yan process $\bar{p}p\to Z_2\to\bar{l}l$ is used as normalization. For numerical estimations, we used the parameters determined by electroweak data analysis [27] with pure vector couplings of Z_2. One would expect the amplitude for the process $\bar{p}p\to Z_2\to\bar{l}lh$ to be enhanced by the factor $\rho_{ht}^2\ V'\ V$ with respect to the standard model prediction. The calculated results for the decay width for $Z_2\to\mu^+\mu^-h$ normalized to $Z_2\to\mu^+\mu^-$ are given in Fig. 8.

We used the typical value for the mixing angle $\theta=-0.02$ for a reasonable approximation $\sin\theta\simeq\theta$ and the trial value $c=0.02$ for the ratio Γ_{Z_2}/M_{Z_2} [27]. Note that the changing of the parameters θ and c in the window of allowed electroweak parameters [27] does not lead to visible new effects with respect to the estimation done in Fig. 8.

VI. CONCLUSION

To summarize, we have demonstrated that the bounds on the scale of quark-lepton compositeness and the Z' boson mass in the extended $SU(2)_h\times SU(2)_l$ model, which are derived from the data taken at the Tevatron (CDF analysis), can be combined to constrain the upper limit of the masses of mass eigenstates $\tilde{t}_1$ and $\tilde{t}_2$ and thus can be used to sensitively probe radiative corrections to the MSSM Higgs sector. Comparison of experimentally measured radiative corrections combined into Δ with calculations can give a precise estimation of the lower bound on the h (as well as A) Higgs boson masses. The analysis of the scale Λ as well as the precise measurement of the lower bound on the Z' boson mass at the Tevatron run II can probe the *CP*-violating mixing between two heavy neutral *CP* eigenstates H and A, and, as a consequence, the nonminimality of the MSSM Higgs sector. It is expected that the Tevatron run II experiments will be able to exclude Z' bosons with masses up to 750 GeV. This leads to a restriction of model scale parameters like $x=u^2/v^2$, which would grow. Note that recent experimental limits on W'_{LR} and Z'_{LR} gauge bosons in the canonical left-right symmetric model [32] require that their mass be higher than about 800 GeV. An important question is whether the forthcoming data at the Tevatron run II at $\sqrt{s}=2$ TeV will progress far enough to determine the lower bounds on the Λ scale and the Z' boson mass within the models considered in this work.

A large decay width $h\to gg$ arises due to a sum of one-loop intermediate heavy quark states largely decaying into two gluons (a new interaction of the h Higgs boson is extended only to the third family). The $h\to\bar{b}b$ channel will be diminished, otherwise it must have already been established at the Tevatron. The $\bar{l}l$ mode via gluon fusion may be significantly enhanced in the MSSM. Thus, this mode is a very promising channel to discover the *CP*-odd Higgs boson A. The discovery region for the $\tau^+\tau^-$ mode might be enlarged as well as the $\mu^+\mu^-$ channel, if a precise reconstruction of the mass of the A Higgs bosons were available. The detection modes $A\to\gamma\gamma h$ and $A\to ggh$ may be good channels for discovery of two Higgs bosons predicted in the MSSM.

From the above considerations, we note that the Tevatron run II might obtain evidence of Higgs bosons (h, H, and/or A) whose dominant final states could be gluon jets and/or pairs of $\tau^+\tau^-$ or $\mu^+\mu^-$ instead of $\bar{b}b$.

ACKNOWLEDGMENTS

We thank G. Arabidze for help in numerical calculations.

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