On search for new Higgs physics in CDF at the Tevatron

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Abstract

We discuss the Higgs boson mass sum rules in the Minimal Supersymmetric Standard Model in order to estimate the upper limits on the masses of stop quarks 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 new extra gauge boson search is 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 the MSSM Higgs bosons in both CP-even and CP-odd sectors and also some extra gauge bosons.

1 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 $G' \subset Z'$, $W^{\pm'}$, $Z''$, $W^{\pm''}$, ... have not yet been established and physics of those particles still remain elusive. In recent years, the amount of works for searching for the Higgs and extra gauge bosons have considerably increased and this research subject is now one of the most exciting topics in searching physics beyond the standard model (SM). On one hand, recent LEP 2 experiments determined a lower bound of the Higgs boson mass to be approximately 114.1 GeV [1]. Furthermore, it should be pointed out that the 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 are translated into lower bounds on the masses of extra neutral gauge bosons $Z'$. On the other hand, recent theoretical progress on study of strong and electroweak interactions and the fundamental concept on natural extensions of the SM, have turned the physics of interplay between the known matter fields and exciting new phenomena related to new matter fields via new interactions.

One of the most challenging current topics beyond the SM is to study physical implication of a set of Higgs particles and extra gauge bosons predicted by the Minimal Supersymmetric Standard Model (MSSM) in the Tevatron $\bar{p}p$ collider experiment at $\sqrt{s}=2$ TeV. The production of Higgs bosons or their decays (via the heavy quark/lepton interactions) are promising and useful processes which would be searched for in the forthcoming high energy experiments. With the advent of the Run II at the Tevatron, it is expected to be clarified that the intermediate heavy quark loop can be a sizeable 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 $G'$, existence of those neutral gauge bosons are required by the 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 $G'$-bosons is rather uncertain, though these bosons could be heavy enough with the 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 the 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 strength and thus could give the answer to the question of why the top quark is so heavy, since they might single out the
fermions of the third generation. The precision measurements of electroweak parameters narrowed the allowed region of extra gauge boson masses keeping Higgs boson masses to be 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, new extra gauge bosons, and to give the estimation of the upper limits on the masses of stop quarks as well as the lower bounds on the masses of the scalar Higgs bosons both in 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 of the quantity $m_{\tilde{t}_1} \cdot m_{\tilde{t}_2}$, i.e. the product of masses of stop eigenstates $\tilde{t}_1$ and $\tilde{t}_2$. We also discuss how the lower bounds of the scalar Higgs boson masses can be obtained from the forthcoming Tevatron data. Furthermore, we emphasize that the forthcoming experiments for discovering the supersymmetry in both Higgs and quark sectors could lead to estimations of the masses of neutral and charged extra gauge bosons $Z'$ and $W'^\pm$, respectively. An example-the models in which the precision electroweak data allow the extra gauge bosons to be of the order $O(0.5 \text{ TeV})$ are, e.g., the noncommuting extended technicolor models [5]. The $Z'$ and $W'^\pm$ 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 like

$$\bar{p}p \rightarrow gg \rightarrow (h/H)X, AX$$

keeping in mind that the main channels in the MSSM are the scalar ones, i.e., $\bar{ll}$, $QQ$, $\gamma\gamma h$, $ggh$ ($\gamma$, $g$, $l$ 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 photons ($\gamma\gamma$) or two gluons ($gg$) decays the invariant mass of $\gamma\gamma$ 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 $h, H \rightarrow \gamma\gamma$ and $h, H \rightarrow gg$ as well. It is known that at high energy $\bar{p}p$ (or $pp$) collisions a contribution of gluonic interactions become large due to increase of gluon densities in the proton. The Tevatron Run II experiment could observe the lightest Higgs boson production via a two gluon fusion $gg \rightarrow h$ with the cross-section $\sigma$ of an order $\sigma \sim O(1.0 \text{ pb})$ at $m_h \sim 110 \text{ GeV}$ [10]. Raise of the $h$-boson mass up
to 180 GeV would lead to decreasing of $\sigma$ up to the order $\mathcal{O}(0.2 \text{ pb})$. For the $A$-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., $A \rightarrow \bar{\ell}\ell, \gamma\gamma h, ggh$ (here, the leptons $\ell$ run over electrons $e$, muons $\mu$ and $\tau$-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

$$m_h^2 + M_H^2 = M_A^2 + m_Z^2$$

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 Section 2, we define the model. Estimation of the upper limits of the masses of stop quarks as well as the lower bounds on the masses of the scalar Higgs bosons will be discussed in Sec. 3. Section 4 is devoted to study on rare decays of $h$-, $H$-, and $A$-Higgs bosons in the MSSM. Section 5 focuses on the $h$ Higgs boson production in the decay of an extra gauge boson $Z_2$. Finally, in Section 6, we give our conclusions.

## 2 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 (labels $h$ and $l$ mean heavy and light, respectively) the gauge boson eigenstates are given by [12]

$$A^\mu = \sin \theta (\cos \phi W_{3h}^\mu + \sin \phi W_{3l}^\mu) + \cos \theta X^\mu$$

for a photon and

$$Z_1^\mu = \cos \theta (\cos \phi W_{3h}^\mu + \sin \phi W_{3l}^\mu) - \sin \theta X^\mu,$$

$$Z_2^\mu = -\sin \phi W_{3h}^\mu + \cos \phi W_{3l}^\mu$$

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

\[
\left( \begin{array}{c}
Z \\
Z'
\end{array} \right) \simeq \left( \begin{array}{c}
\frac{1}{\cos \phi} \frac{\sin \phi}{x \cos \theta} \\
-\frac{\cos \phi \sin \phi}{x \cos \theta} \frac{1}{1}
\end{array} \right) \left( \begin{array}{c}
Z_1 \\
Z_2
\end{array} \right),
\]

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

The parameter $x$ in (9) 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 $v \approx 246$ GeV is the vacuum expectation value of the (composite) scalar field responsible for the symmetry breaking $SU(2)_L \times U(1)_Y \rightarrow U(1)_{em}$ in the model of extended weak interactions. The generator of the $U(1)_{em}$ group is the usual electric charge operator $Q = T_3^h + T_3^l$.

At large values of $\sin \phi$, the $Z_2$-boson has an enhanced coupling to the third generation fermions through the covariant derivative

\[
D^\mu = \partial^\mu - i \frac{g}{\cos \theta} Z_i^\mu \left( T_{3h}^i + T_{3l}^i - \sin^2 \theta \cdot Q \right)
\]

\[-i g Z_2^\mu \left( -\frac{\sin \phi}{\cos \phi} T_{3h}^i + \frac{\cos \phi}{\sin \phi} T_{3l}^i \right). \]

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

\[
\mathcal{L} \supset \frac{1}{\Lambda_{LL}^2} \left[ g_0^2 (\bar{E}_L \gamma_\mu E_L)(\bar{Q}_L \gamma_\mu Q_L) + g_1^2 (\bar{E}_L \gamma_\mu \tau_a E_L)(\bar{Q}_L \gamma_\mu \tau_a Q_L) \right]
\]

\[+ \frac{g_e^2}{\Lambda_{LR}^2} (\bar{e}_R \gamma_\mu e_R)(\bar{Q}_L \gamma_\mu Q_L) \]

\[+ \frac{1}{\Lambda_{LR}^2} (\bar{E}_L \gamma_\mu E_L) \frac{1}{\Lambda_{RR}^2} (\bar{e}_R \gamma_\mu e_R) \sum_{q=u,d} g_q^2 (\bar{q}_R \gamma_\mu q_R), \]

(8)

where $E_L = (\nu_e, e)_L$, $Q_L = (u, d)_L$; $g_i$ are the effective couplings and $\Lambda_{ij}$ are the scales of new physics. The aim of the CDF collaboration analysis [2] was to search for the deviation of the SM prediction in the dilepton production spectrum. If no such deviations have been found, the lower bound of the $\Lambda$-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 to the following part of the Lagrangian density (see [9])

\[
\mathcal{L} \supset -\frac{g^2}{M_2^2} \left( \frac{\cot \phi}{2} \right)^2 \left( \sum_{l=e,\mu} \bar{l}_L \gamma_\mu l_L \right) \left( \sum_{q,u,d,s,c} \bar{q}_L \gamma_\mu q_L \right) , \]

(9)
where \( g = e / \sin \theta \). We suppose that the couplings in the first two generations are same in strength.

### 3 Mass bound on stop quarks 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 \( Z' \)-boson; \( \delta_{ZZ'} = M_{Z'}^2 - m_Z^2 \). The correction \( \Delta \) reflects the contribution from loop diagrams involving all the particles that couple to 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, \( m_t \) and \( m_W \) are the masses of top quark and \( W \)-boson, respectively. \( \tan \beta \) defines the structure of the MSSM. The values of \( \Delta \sim \mathcal{O}(0.01 \ TeV^2) \) 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 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 parameters space and the experimental data \(^{[17]}\).

Comparing \(^{(8)}\) and \(^{(9)}\), one can get the following relation between \( M_{Z'} \) and \( \Lambda \) as \(^{[9]}\)

\[
M_{Z'} = \sqrt{\alpha_{em}} \Lambda \cot \phi/(2 \sin \theta) , \tag{12}
\]

where the value of \( \Lambda \) 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 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 $\alpha$ 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 (10). Here, $\mu$ is the positive massive parameter which can, in principle, be defined from the 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.

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 (10). We suppose that $M_{Z}'$ in the l.h.s. of (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 r.h.s. of (10) is model-dependent, where, to the 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 $\Lambda$ [2] entering into (12), one can easily find the upper limit of the product of $m_{t_1} \cdot m_{t_2}$ from the following relation [13]

$$\Delta < (B + M_{H}^*)(B - f \; C) + m_{H}^2 - m_{Z}^2 (1 - \sin^2 2\beta) + \mu^2,$$

where $M_{H}^* = (M_{A}^2 + m_{Z}^2 \sin^2 2\beta + \mu^2)^{1/2}$, $f \equiv f(\phi) = \cot \phi \sqrt{\alpha_{em}}/(2 \sin \theta)$, $B \equiv B(x, \phi) = m_{W} \sqrt{x}/(\cos \phi \sin \phi)$, and $C$ is a minimal value of the $\Lambda$ 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 $\Delta$ in (11) and the mass $M_{H}$ defined in the decoupling regime. Thus, the upper limit on $m_{t_1} \cdot m_{t_2}$ 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. 1 shows the upper limit on $L \equiv \log(\frac{m_{t_1} \cdot m_{t_2}}{m_{t_2}^2})$ as a function of sin $\phi$ for $x = 2$ and $x = 3$ at fixed values of $\mu$ and $M_{A}$. We use the following range of the mixing parameter $0.75 \leq \sin \phi \leq 0.85$ 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 $3.9 \geq x \geq 1.6$ (left-handed muons and up-type quarks is 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 $A$-Higgs boson are $M_A > 88.4$ GeV \cite{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 $g/g_t$. 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 $g_h$ is originated 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 the Fig. 1, the width $\Gamma_{Z'}$ of the $Z'$-boson falls to a minimum in the neighborhood of $\sin \phi = 0.8$ \cite{9}, due to the decreasing couplings to 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.

Fig. 1 The upper limit on $L \equiv \log \left( \frac{m_{t1} m_{t2}}{m_{t3}^2} \right)$ 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; $\tan \beta = 30$. 
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.

The CDF analysis of the contact interaction between left-handed muons and the 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 $p\bar{p} \rightarrow AZ + X$ with $\sqrt{s} > M_A + m_Z$. Here, we did not use the mass difference between $t_1$ and $t_2$ 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 a 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]:

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

The analytical expressions for the loop corrections are unknown yet. However, it is known that the one-loop corrections can be summarized into the 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 $\sin \phi$:

$$\sum_{j=1}^2 M_{H_j}^2 > M_A^2 + m_Z^2 - m_h^2 + \frac{\alpha_m C^2 \cot^2 \phi}{4 \sin^2 \theta}.$$  \hfill (15)

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

![Fig. 3](image)

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

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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 electron (muon) and up-type quarks: $\Lambda > 4.1$ TeV and $\Lambda > 3.7$ TeV for left-handed muons and left-handed electrons, respectively.

4 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 on the decays $X \to gg$ and $X \to \gamma\gamma$ ($X : h, H$), where the former dominates over the later. The decay width of $X \to gg$ (the radiative corrections are not included) is given by

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

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

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

Here, $\tau_Q = (2m_Q/m_X)^2$, $\xi_Q = \sqrt{1 - \tau_Q}$, $\rho_{hQ} = -\sin \alpha/\cos \beta$ 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 $\alpha$ diagonalizes the CP-even Higgs squared-mass matrix. For comparison, the decay width $\Gamma(X \to \gamma\gamma)$ can be easily obtained from the corresponding decay to two gluons [10] by a simple replacement $\alpha_s \to \frac{3}{\sqrt{2}} \alpha e_Q^2$ (the heavy quark charges $e_Q$ are included into the formfactors) and $\tilde{F}_Q \to F_Q + F_W + F_H^\pm$

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

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where the formfactor $F_Q$ contributed from heavy quarks $Q$ is given by

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

(19)

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

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

(20)

and

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

(21)

respectively. Here, $\rho_{hW} = \sin(\beta - \alpha)$, $\rho_{HW} = \cos(\beta - \alpha)$, $\rho_{hH^\pm} = \rho_{hW} + \frac{\cos 2\beta \sin(\beta + \alpha)}{2 \cos^2 \theta_W}$, $\tau_{H^\pm} = (2M_{H^\pm}/m_X)^2$. Calculated results of $\Gamma(H \rightarrow gg)$, $\Gamma(h \rightarrow gg)$ and $\Gamma(H \rightarrow \gamma\gamma)$, $\Gamma(h \rightarrow \gamma\gamma)$ are given in the 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 and left-and right-handed sfermions as well as $h$-Higgs to charginos were neglected. The following relation between the mixing angle $\alpha$ and $\tan \beta$

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

(22)
Fig. 4 The decay width $\Gamma(H \rightarrow gg)$ as the function of $M_H$ for different $\alpha_s$ (from below $\alpha_s = 0.110$, $0.115$, $0.119$, $0.123$) and $\tan \beta$ = 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 $gg$ and $\gamma\gamma$ channels or $\bar{b}b$ final state.

(b) The decays $A \rightarrow \bar{\ell}\ell$.

In the MSSM, the decay amplitude of a CP-odd neutral Higgs boson $A$ (with mass $M_A$ and 4-momentum $P_{\mu}$) into a lepton $\ell$ and antilepton $\bar{\ell}$ with 4-momenta $p_{\mu}$ and $(P - p)_{\mu}$, respectively, is given by the following Feynman amplitude

$$A m(A \rightarrow \bar{\ell}\ell) = i \sum_{Q,b,t} \bar{u}(p) \gamma_5 F_Q(M_A^2) v(P - p) ,$$

where $F_Q(t = M_A^2)$ is the complex function (formfactor) providing the transition from the $A$ Higgs boson into a lepton pair. One can consider this tran-
Fig. 5 The decay width $\Gamma(h \to gg)$ as the function of $m_h$ for different $\alpha_s$ (from below $\alpha_s = 0.110, 0.115, 0.119, 0.123$) in the decoupling limit.

The transition via the intermediate of-shell spin-1 bosons $V^*$ (photons, $W^\pm$-bosons, and so on) within the vertex $AV^*V^*$

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

with

$$
F_A[k^2, (P - k)^2] = \sum_{Q:b,t} \rho_Q \epsilon_Q^2 f_{VV} F_{QA}[k^2, (P - k)^2].
$$

Here, $\rho_Q = \tan \beta$ for $b$-quark-loop and $\rho_Q = \cot \beta$ for $t$-quark-loop; $k_\mu$ is the 4-momentum of a $V^*$-boson. There is a natural normalization condition $F_A(0,0) = f_{VV}$ in (24) where the decay constant $f_{VV}$ is given for the decay of $A$-Higgs into two real $V$-bosons. The form of $F_A$ in (24) should be taken so that the Feynman integrals are not divergent. Here, we simply take $F_{QA}[k^2, (P - k)^2]$ as

$$
F_{QA}[k^2, (P - k)^2] = \frac{\lambda_Q^2}{\lambda_Q^2 - k^2 - (P - k)^2}
$$
which has the good analytical properties, and carries both the vector dominance features and the properties of the right static limit. The parameter $\lambda_Q$ 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 $A \rightarrow \bar{ll}$, one has to take into account the fact that the formfactor $F_A(k^2, (P - k)^2)$ should be complex function with the following real part

$$Re \ F_A(M_A^2) = \frac{P}{\pi} \int_{0}^{\infty} \frac{Abs[F_A(t)]}{t - M_A^2} \ dt . \quad (27)$$

Here, the absorptive part $Abs[F_A(t)]$ can be found if we take the $S_{10}^1$ lepton pair in the final state and taking into account the unitarity condition for the decay amplitude

$$Abs[F_A(t)] = \frac{1}{2i \sqrt{2} t} \int d\Omega \ \delta^4(P - \sum_l q_l) \langle I | T | \bar{ll} \rangle_{S_{10}^1}^* \ Am(A \rightarrow I) . \quad (28)$$
Fig. 7 The decay width $\Gamma(h \rightarrow \gamma\gamma)$ as the function of $m_h$ in the decoupling limit.

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

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

with

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

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 cancelled for the $VV$-channel. The corresponding SM background comes from the Drell-Yan production

$$\bar{q}q \rightarrow \gamma^*, Z^*, Z^{**} \rightarrow \bar{l}l.$$
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 1 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 \( M_A \).

Table 1. The values \( \text{Br}(A \to \bar{u}/\gamma \gamma) \) as the function of \( M_A \) where \( l = \mu^- \) and \( \tau^- \).

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

At the end of this part, we conclude that apart from reducing of the \( A \to \tau^+ \tau^- \)-signal due to some experimental constraints, this decay mode could be detected at the Tevatron’s Run II as (almost) 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 to investigate 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 can be more promising for discovery for \( A \)-Higgs bosons than that of \( \tau^+ \tau^- \) or \( \mu^+ \mu^- \) channels.

The matrix element of the decay of \( A \)-Higgs bosons with momentum \( P \) into \( \gamma, \gamma \) and \( h \) with the momenta \( k_1, k_2 \) and \( k_3 \), respectively, has the following form

\[
M(A \to \gamma \gamma h) = \frac{\sqrt{N_c} \pi \alpha i g}{m_W \sqrt{M_A}} \text{Tr} \left[ \sum_{Q,b,t} \gamma_5 \rho_Q c_Q^2 m_Q (M_A - \hat{P}) \Gamma_Q \rho_{hQ} T_Q \right],
\]

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); \( \Gamma_Q \) is the vertex function describing the couplings \( h \bar{Q}Q \) taking into account the \( \rho_Q \) flavor-dependent factor. The differential distribution of the decay width \( A \to \gamma \gamma h \) over the invariant mass \( \hat{s} \) of final state’s two-photon pairs and normalized to the \( \gamma \gamma \) width is given by

\[
\frac{1}{\Gamma(A \to \gamma \gamma)} \frac{d\Gamma(A \to \gamma \gamma h)}{d\hat{s}} = \rho_{hQ}^2 \frac{G_F a^2 \hat{s} M_A^2}{8 \sqrt{2} \pi^2 (a^2 - \hat{s}^2)^2 (a + \hat{s}/2)},
\]
where \( a = (1 - \kappa^2)/2 \), \( \kappa = m_h/M_A \), \( \bar{s} = s/M_A^2 \), \( s = (k_1 + k_2)^2 \). In Table 2, we present \( Br(A \rightarrow \gamma\gamma h/\gamma\gamma) \) as a function of \( \kappa \) for different values of \( \tan \beta = 5, 10, 20 \) and 50 at fixed \( M_A=0.8 \) TeV.

**Table 2.** \( Br(A \rightarrow \gamma\gamma h/\gamma\gamma) \cdot 10^{-2} \) as a function of \( \kappa \) for \( \tan \beta = 5, 10, 20 \) and 50; \( M_A \) is taken to be \( M_A=0.8 \) TeV as a typical value.

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

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

Compared with \( A \rightarrow \tau^+\tau^-/\mu^+\mu^- \) cases as discussed in (b), we conclude that the \( A \rightarrow \gamma\gamma h \) channel in MSSM could be observed for \( 0.15 \leq \kappa \leq 0.2 \) \( (120 \text{ GeV} \leq m_h \leq 160 \text{ GeV}) \) and for \( \tan \beta > 5 \).

## 5 The decays \( Z_2 \rightarrow Z_1 h \)

To reach a unified theory of all interactions one could start with 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 \( W^{\pm'} \) ones. Experimental signals of those extra gauge bosons would give very important information about the underlying GUT and its origin. Search for \( Z' \) and \( W^{\pm'} \) 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 \( W^{\pm'} \) because they are sensitive to rare processes including their decays.

The cross-section \( \sigma(\bar{p}p \rightarrow Z'\bar{l}l) \) for \( Z' \) production at the Tevatron is inversely proportional to the total decay width \( \Gamma_{Z'} \). Obviously, the inclusion of extra channels to \( Z' \)-decays leads to that \( \Gamma_{Z'} \) will become larger and
σ(¯¯pp → Z′ll) smaller. As already pointed out [25,26], if the decays of Z′ into
h + {QQ} s=1, hll are kinematically allowed, search for h Higgs boson or an
exotic heavy quark-antiquark bound states {QQ} s=1 with spin s = 1 would
be the interesting channels. The study of the decay Z′ → h + {QQ} s=1 [25,26]
would provide a useful information about the nature of the extended gauge
structure such as the couplings of Z′ with heavy quarks in both vector and
axial-vector sectors, couplings of h with heavy quarks, Z − Z′ mixing effects,
physics of mass eigenstates Z1 and Z2 and interplay with matter fields.

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

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

where \(\Psi_f\) is the fermion field with the flavor \(f\). The first term in (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, \(T_3^L\) 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
\((g_v^{f'}, g_a^{f'})\) represent the chiral properties of interactions of Z- and Z'-bosons
to \(\Psi_f\), respectively. The mass eigenstates Z1 and Z2 are parameterized by a
mixing angle \(\theta\) originated 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

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

where

\[
V_f = g_v^f \cos \theta + \frac{g_2}{g_1} g_a^{f'} \sin \theta, \quad A_f = g_a^f \cos \theta + \frac{g_2}{g_1} g_v^{f'} \sin \theta \quad (36)
\]

and

\[
V^{f'} = \frac{g_2}{g_1} g_v^{f'} \cos \theta - g_v^f \sin \theta, \quad A^{f'} = \frac{g_2}{g_1} g_a^{f'} \cos \theta - g_a^f \sin \theta \quad (37)
\]
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 \( \sin^2 \theta_W(M_Z) = 0.23117 \) [18].

The decays of \( Z_2 \) are the promising place to search for CP-even light Higgs-boson \( h \). Here, the effects of heavy quarks cannot be neglected. As a result, there exist 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 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 interactions of \( h \) is reproduced by rescaling all the mass terms \( m_j = m_j (1 + h/v), j= \) quarks, W, Z, Z',..., and \( \alpha_s \rightarrow \alpha_s + \delta \alpha_s \) with \( \delta \alpha_s = \alpha_s^2 h/(3 \pi v) \) [28-30]. Here, the number of heavy quarks is restricted by one, that is only top quark loop will be involved into the game. The interaction Lagrangian looks like:

\[
\mathcal{L}_{\text{int}} = (1 + h/v) \rho_{ht} m_t \bar{t} t + \frac{\alpha_s \rho_{ht}}{12 \pi v} h G_{\mu\nu} G^{\mu\nu} - (1 + h/v)^2 \left( m_W^2 W_\mu^+ W^\mu + \frac{1}{2} m_Z^2 Z_\mu^+ Z_\mu + \frac{1}{2} m_{Z'}^2 Z'_\mu^+ Z'_\mu \right),
\]

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 \( h \) Higgs boson is low enough for the decay \( Z_2 \rightarrow Z_1 h \rightarrow \bar{ll}h \) so as 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 \( \Gamma(Z_2 \rightarrow \bar{ll}h) \) over the dimensionless variable \( x = (p_l + p_{\bar{t}})^2/M_{Z_2}^2 \) (\( p_l \) and \( p_{\bar{t}} \) are the momenta of a lepton and antilepton, respectively) is given by

\[
R_{\bar{ll}h}(x) = \frac{1}{\Gamma(Z_2 \rightarrow \bar{ll})} \frac{d \Gamma(Z_2 \rightarrow \bar{ll}h)}{dx} = \frac{g^2 V' V M_{Z_2}^2 \rho_{ht}^2}{4 \cos^2 \theta_W V' 24 \pi^2 v^2 x} \left[ (1 - a_h)^2 + x^2 - 2 x (1 + a_h) \right]^{1/2} \times \left( 1 - \frac{4}{x} a_l \right)^{1/2} \left( 1 + \frac{2}{x} a_l \right) \frac{(1 - a_h)^2 + x^2 + 2 x (2 - a_h)}{(1 - x)^2 + c^2},
\]

where \( V' = 0.62 \lambda^{1/2} - \theta g_v, V = g_v + 0.62 \lambda^{1/2} \theta \) [27], \( V' = (\sqrt{3}/2) (1 - 4 \sin^2 \theta_W)^{1/2} \) [31], \( a_j = (m_j/M_{Z_2})^2 \) with \( j= l, h; c = \Gamma_{Z_2}/M_{Z_2} \). To get \( V' \) 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 (34), the Drell-Yan process $\bar{p}p \rightarrow Z_2 \rightarrow \bar{\ell}\ell$ 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 $\bar{p}p \rightarrow Z_2 \rightarrow \bar{\ell}\ell h$ to be enhanced by the factor $\rho_h^2 V'V$ with respect to the standard model prediction. The calculated results of the decay width for $Z_2 \rightarrow \mu^+\mu^- h$ calculations normalized to $Z_2 \rightarrow \mu^+\mu^-$ are given in Fig. 8.

![Image](image.png)

Fig. 8 The $x = (p_\perp + p_{\ell \perp})^2/M_{Z_2}^2$ distribution of $R_{\mu^+\mu^- h}$ in the decay $Z_2 \rightarrow \mu^+\mu^- h$ for $M_{Z_2} = 0.3$ TeV, 0.5 TeV, 0.8 TeV and 1 TeV.

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

6 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 $\Delta$ with its 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 $\Lambda$ 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 the restriction of the model scale parameter like $x = u^2/v^2$ which would grow. Note that recent experimental limits on $W^\prime_{LR}$ and $Z^\prime_{LR}$ 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 $A$-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 on the third family). The $h \to \bar{b}b$ channel will be diminished, otherwise it must have already been established at the Tevatron. The $\bar{u}l$ mode via a gluon fusion may be significantly enhanced in MSSM. Thus, this mode is a very promising channel to discover CP-odd Higgs boson $A$. 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 is 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 consideration, we note that the Tevatron Run II can catch an 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 $\bar{b}b$.

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