Enhancement of the gluonic couplings of the Standard Model Higgs Boson through mixing

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Abstract

The mixing of the Standard Model Higgs boson with a \( t\bar{t} \) scalar bound state (arising from a new stronger-than-QCD chiral interaction) can enhance the Higgs boson’s coupling to two gluons. It is possible that the two-gluon decay mode of a light Higgs boson may even dominate. This alters the Higgs boson search strategy at Tevatron Run 2 in a significant way. We present a simple model in which such a scenario may be realized. The modifications to the Higgs search signal, as well as further experimental tests of this scenario are also discussed.

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I. INTRODUCTION

One of the major objectives of Run 2 at the Tevatron is to search for the Higgs boson in the mass range of 110 GeV to 180 GeV \[1\]. Direct searches at LEP2 have shown no clear evidence of the Higgs boson, and have set a lower limit of 114.1 GeV for the Standard Model (SM) Higgs boson \[2\]. Global analyses of LEP1 electroweak data have considerably narrowed down the SM Higgs boson mass \[3\], and prefer a low mass Higgs boson. At Tevatron energy \((\sqrt{s} = 2 \text{ TeV})\), the dominant mechanism of SM Higgs boson production is \(gg \rightarrow h\) with a cross section ranging from 0.9 pb for \(m_h = 110 \text{ GeV}\) to 0.2 pb for \(m_h = 180 \text{ GeV}\). For \(m_h < 135 \text{ GeV}\), \(h\) decays dominantly to \(b\bar{b}\) and this mode of production and decay is hopeless due to the overwhelming QCD background. Form \(m_h > 135 \text{ GeV}\), the \(h \rightarrow WW^*\) decay mode becomes dominant, and the detection of the Higgs boson is promising via this mode. The detection via \(\tau^+\tau^-\) decay mode has not been thoroughly studied and may be on the borderline \[4\]. In the usual scenario, the best way to observe the signal of the SM Higgs boson for \(m_h < 135 \text{ GeV}\) is the production of \(Wh\) via \(q\bar{q}'\) and the subsequent decay \(h \rightarrow b\bar{b}\), where the background can be controlled using appropriate cuts. However, if \(h \rightarrow b\bar{b}\) is not the dominant decay (as in the scenario we shall propose), this signal may be significantly lowered, making the detection in this channel very difficult. On the other hand, if the \(ggh\) coupling is significantly enhanced, Higgs boson production through gluon fusion may be large enough to detect a signal via the \(\tau^+\tau^-\) decay mode for the low mass region \((m_h < 135 \text{ GeV})\), or the signal in the \(WW^*\) decay mode could be large enough to detect in the higher mass range \((m_h > 135 \text{ GeV})\). The object of this work is to explore an alternative scenario in which the \(ggh\) coupling is significantly enhanced, and how that alters the Higgs boson search strategy at Tevatron Run 2 due to the changes of the relative importance of its various production and decay modes.

In the Minimal SM the Higgs boson has a small coupling to two gluons, generated only at one-loop level, a fact that makes its discovery difficult on hadron colliders. In extensions to the SM various Higgs-like scalars are predicted, in some cases with enhanced gluonic couplings.\(^1\) These extensions are pursued in order to have specific explanations for EW

\(^1\) Composite Higgs bosons can couple to gluons directly through their quark constituents. The Higgs mass is typically larger in these models than the range considered in the present paper. See \[5\], however, for an example where \(m_h \leq 330 \text{ GeV}\).
FIG. 1: The exchange of a new gauge boson would generate a large contribution to the $gt\bar{t}$ coupling unless suppressed by $m_{new} \gg 100\, GeV$.

breaking. In this paper we point out that appropriate new physics can generate an enhancement to this coupling even if mechanism of electroweak breaking is left intact. The additional interactions present in such models are not required in order to explain EW breaking, but they cannot be excluded either. Such scenarios must be taken into consideration because they have important phenomenological consequences.

Such an enhancement arises if the Higgs boson mixes with a $t\bar{t}$ bound state. Such a state is not formed in the minimal SM because the top quarks would decay before the strong interaction could bind it\(^2\) (i.e. $\alpha_s(m_t) \sim 0.12$ is too weak). However, the right handed top quark may participate in additional, chiral, new strong interactions. Such interactions are not phenomenologically excluded if the new gauge symmetry is broken and the new gauge boson picks up a mass $\mathcal{O}(m_{new})$.\(^3\) We denote by $\sigma$ a hypothetical $t\bar{t}$ bound state that mixes with the Higgs boson. In spectral notation referring to the $t\bar{t}$ system this state must be a $2^3P_0$ state (to ensure even parity and zero spin.) Such a bound state must be held together by new interactions, whose strength, $\alpha_{new}$, must be much larger than $\alpha_s \sim 0.12$ to form a bound state at all. They must act only on the third family (otherwise they would have already been detected in processes such as $q\bar{q} \to b\bar{b}$ at the Tevatron.) The dangerous place is the $gt\bar{t}$ vertex, where one loop corrections, due to the exchange of the new gauge

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\(^2\) The effects of QCD interactions on a $t\bar{t}$ system are small. Approximating the QCD effects by a Coulomb potential, we would find a bound state whose revolution time is $\tau_r = \frac{9}{4m_t\alpha_s^2} \sim 1\, GeV^{-1}$. Because the top quarks constituting the bound state have a decay width $\Gamma_t = 1.2\, GeV$, the would-be bound state decays before even one revolution.

\(^3\) An interaction in which left handed quarks participate is excluded by their contribution $Z \to b\bar{b}$, provided $SU(2)_W$ is respected. Anomaly cancellation might be arranged through a new fermion sector which may be too heavy to be presently observed.
bosons might have a large effect, as shown in Fig. 1. At the Tevatron, only the $t\bar{t}$ production cross section will be affected. This cross section is known with a $\sim 10\%$ precision, and the effect is $\mathcal{O} \left( \frac{\alpha_{\text{new}}}{4\pi} \times \frac{m_t^2}{m_{\text{new}}^2} \right)$, so we would prefer a new interaction effective at scales much above $\sim 100 \text{ GeV}$.

In this paper our goal is not to build a specific model. Nevertheless we point out that the new interactions, being chiral, will be very different from a scaled-up copy of QCD. A consequence of this fact is that the standard way of calculating quarkonium decays, Nonrelativistic QCD [8], is not applicable here (the size of the NRQCD operators are determined by QCD dynamics.) We’ll do a potential model calculation instead. We use a Yukawa potential to model the $t\bar{t}$ interaction and argue that the strength of the Yukawa potential will be related to $\alpha_{\text{new}}$, and that the screening length $m^{-1}$ is related to $m_{\text{new}}^{-1}$. These parameters, however, may not exactly coincide, leaving us room to argue that the screening length $m^{-1}$ in the potential may be somewhat longer than the inverse mass $m_{\text{new}}^{-1}$ suppressing visible effects in top production at the Tevatron.

In Sec. II we work out the general formalism of the mixing between a light Higgs $h$ and the $\sigma$, and show that, with the exception of a small part of the parameter space, the gluonic branching ratio of the Higgs is determined directly by the gluonic branching ratio of the $\sigma$ and also by the masses of the $h$ and $\sigma$ (and is independent of any mixing parameters.) Then we proceed in Sec. III to calculate the binding energy and wave functions of the lowest lying states. We find that, in order for a $2^3P_0$ state to form and not to have a too wide $\sigma \to b\bar{b}$ width, the interaction strength is restricted to be between $0.5 \leq \alpha_{\text{new}} \leq 1.0$, while the screening mass can still exceed $m > 100 \text{ GeV}$. However, in such a case the binding of the $1S$ states is so deep that their masses become negative. The solution is spin-dependent potentials. Nuclear physics is witness that strong interactions, even nonchiral, may generate very strong $L \cdot S$ couplings. A spin-orbit interaction that binds states with negative values of $L \cdot S$ stronger, can lead to a situation when the only bound state that is formed is $\sigma \equiv 2^3P_0$, as we show in a numerical example. We do not think that a one-loop calculation of the relative strengths of the spin independent and the $L \cdot S$ terms in the potential is reliable, so that we handle this ratio as an unknown input parameter.^[4]

[^4]: We also note that higher order (in $\alpha_{\text{new}}$) effects and/or relativistic corrections may also induce a mixing between the Higgs boson and the $1^1S_0$ state. At tree level such mixing is absent and the interaction is repulsive in the singlet and attractive in the triplet channel which does not allow the formation of a $1^1S_0$. 

In Sec. [IV] we argue that the presence of such mixing can lead to the enhancement of final states with two gluons from Higgs decay in lepton colliders. A large $h \rightarrow gg$ branching ratio can lead to a decrease of the $h \rightarrow b\bar{b}$ signal, dominant in the Minimal SM, and this can adversely affect Higgs searches in $e^+e^-$ colliders. On hadron colliders the increased gluonic coupling enhances Higgs production in gluon fusion. At the Tevatron, perhaps the most interesting enhancement occurs in $gg \rightarrow h \rightarrow \tau^+\tau^-$ for a light Higgs ($m_h < 135 \text{ GeV}$), while heavy Higgs production ($m_h > 135 \text{ GeV}$) will be enhanced in $gg \rightarrow h \rightarrow W^+W^−, WW^*$. 

II. MIXING AT THE LAGRANGIAN LEVEL

The most general form of the relevant part of the Lagrangian after the new interactions have been integrated out is

$$L_{mix} = \frac{1}{2} \partial_\mu \Phi^\dagger_0 \cdot S \cdot \partial^\mu \Phi_0 - \frac{1}{2} \Phi^\dagger_0 \cdot M^2_0 \cdot \Phi_0$$

with \( \Phi_0 = \begin{pmatrix} h_0 \\ \sigma_0 \end{pmatrix} \), and \( S, M \) representing real symmetric \( 2 \times 2 \) mixing matrices. This can always be diagonalized by a field redefinition \( \Phi = A^T \times \Phi_0 \) satisfying \( A \cdot S \cdot A^T = 1 \), resulting in

$$L_{mix} = \left[ \frac{1}{2} (\partial h)^2 - \frac{m^2_h}{2} h^2 \right] + \left[ \frac{1}{2} (\partial \sigma)^2 - \frac{m^2_\sigma}{2} \sigma^2 \right] + h j_h + \sigma j_\sigma. \quad (2)$$

The mixing can be parametrized by four parameters, \( \vartheta, \theta, \rho_h \) and \( \rho_\sigma \),

$$A = \begin{bmatrix} \rho_h \cos \theta & -\rho_\sigma \sin \vartheta \\ \rho_h \sin \theta & \rho_\sigma \cos \vartheta \end{bmatrix}. \quad (3)$$

Note that the kinetic mixing with these parameters is

$$S = (A^T \cdot A)^{-1} = \frac{1}{\cos^2(\theta - \vartheta)} \begin{bmatrix} 1/\rho^2_h & \sin^2(\theta - \vartheta)/\rho_h\rho_\sigma \\ \sin^2(\theta - \vartheta)/\rho_h\rho_\sigma & 1/\rho^2_\sigma \end{bmatrix}, \quad (4)$$

so that in the absence of kinetic mixing the parameters reduce to only one mixing angle through \( \theta = \vartheta \) and \( \rho_h = \rho_\sigma = 1 \).

The interactions of the physical Higgs and \( \sigma \) are described by the currents in (2),

$$j_h = \rho_h \cos \theta j_{h_0} - \rho_\sigma \sin \vartheta j_{\sigma_0} \quad \text{and} \quad j_\sigma = \rho_\sigma \cos \vartheta j_{h_0} + \rho_h \sin \theta j_{\sigma_0},$$

state. This conclusion may be incorrect though when higher orders are incorporated.
which in the absence of kinetic mixing reduces to

\[ j_h = \cos \vartheta j_{h_0} - \sin \vartheta j_{\sigma_0} \quad \text{and} \quad j_\sigma = \cos \vartheta j_{h_0} + \sin \vartheta j_{\sigma_0}. \] (6)

The effects of such mixing will show up both in \( h \) production and decay.

In the following we will focus on the case of a light Higgs with mass \( 100 \text{GeV} \leq m_h \leq 135 \text{GeV} \), when only the \( g, g \) and \( b \bar{b} \) Higgs decay channels are important in the SM. We will comment on the heavy Higgs in Sec. IV. The original currents are

\[ j_{h_0} = y_b \overline{b^4} b_A, \quad j_{\sigma_0} = \frac{f_g}{2m_t} \times G_a^{\mu\nu}G_\mu^\nu + \bar{b} \left[ f_R \frac{1 + \gamma^5}{2} + f_L \frac{1 - \gamma^5}{2} \right] b. \] (7)

These depend, in general, on three new dimensionless couplings \( f_R, f_L \) and \( f_g \). In the above the \( h_0gg \) coupling has been neglected, because it only arises from loop diagrams.

After diagonalization the currents that couple to the physical Higgs and \( \sigma \) become

\[ j_h = \bar{b} \left[ (y_b \cos \theta - f_R \sin \vartheta) \frac{1 + \gamma^5}{2} + (y_b \cos \theta - f_L \sin \vartheta) \frac{1 - \gamma^5}{2} \right] b - \frac{f_g \sin \vartheta}{2m_t} \times G_a^{\mu\nu}G_\mu^\nu, \] (8)

\[ j_\sigma = \bar{b} \left[ (y_b \sin \theta + f_R \cos \vartheta) \frac{1 + \gamma^5}{2} + (y_b \sin \theta + f_L \cos \vartheta) \frac{1 - \gamma^5}{2} \right] b + \frac{f_g \cos \vartheta}{2m_t} \times G_a^{\mu\nu}G_\mu^\nu. \]

The partial decay widths of the Higgs and the \( \sigma \) are found from these couplings at the corresponding poles.

We calculated the total widths of both particles, denoted as \( \Phi = h, \sigma \), coupled to a generic current \( j_\Phi = \frac{f_\Phi}{2m_t} G^2 + \bar{b} \left( f_R^{\Phi \frac{1 + \gamma^5}{2}} + f_L^{\Phi \frac{1 - \gamma^5}{2}} \right) b \), and found

\[ \Gamma_{\Phi \to gg} = \frac{m_\Phi}{32\pi} \left( \frac{m_\Phi}{m_t} \right)^2 \times \left( N_c^2 - 1 \right) \times \left| f_\Phi \right|^2, \]

\[ \Gamma_{\Phi \to b\bar{b}} = \frac{m_\Phi}{32\pi} \times N_c \times \left( \left| f_\Phi \right|^2 + \left| f_{\Phi} \right|^2 \right). \] (9)

The Higgs decay ratio is then found to be

\[ R_h \equiv \frac{\Gamma_{h \to gg}}{\Gamma_{h \to b\bar{b}}} = \frac{N_c^2 - 1}{N_c} \cdot \left( \frac{m_h}{m_t} \right)^2 \times \frac{f_g^2 \sin^2 \vartheta}{y_b^2 \cos^2 \theta + \left( f_L^2 + f_R^2 \right) \sin^2 \vartheta - 2y_b (f_L + f_R) \cos \theta \sin \vartheta}. \] (10)

We note that this formula involves essentially only two parameters. This is manifest if the parameter \( f^2 \equiv f_L^2 + f_R^2 = \frac{32\pi}{N_c m_\sigma} \Gamma_{\sigma_0 \to b\bar{b}} \) is used, together with the angle \( \alpha \), defined as \( f_L = f \cos \alpha, \ f_R = f \sin \alpha \):

\[ R_h = \frac{8 m_h^2}{3 m_t^2} \frac{f_g^2 \sin^2 \vartheta}{(y_b \cos \theta)^2 + (f \sin \vartheta)^2 - 2(y_b \cos \theta)(f \sin \vartheta) \times \sqrt{2} \cos(\pi/4 - \alpha)}. \]
\[
\frac{m_{\sigma}^2}{m_{h}^2} \cdot R_{\sigma} \left( 1 + \left( \frac{y_b \cos \theta}{f \sin \vartheta} \right)^2 - 2 \left( \frac{y_b \cos \theta}{f \sin \vartheta} \right) \times \sqrt{2} \cos \left( \frac{\pi}{4} - \alpha \right) \right),
\]

where we introduced the notation \( R_{\sigma} \equiv \frac{\Gamma_{\sigma \to gg}}{\Gamma_{\sigma \to b\bar{b}}} \). From Eqn. (11) we see that the term containing the parameter \( \alpha \) plays only a minor role (disregarding the exceptional case when a cancellation occurs in the denominator), and this fact justifies the approximation in the last line.

We can argue that except for two very unlikely cases the unity in the denominator of Eqn. (11) dominates. For these exceptions to happen we need \( \frac{\sin \vartheta}{\cos \vartheta} \leq \frac{y_b}{f} \equiv \sqrt{8 \times 10^{4} \Gamma_{\sigma \to b\bar{b}}} \). This would require either (i) a very small mixing angle \( \theta < 10^{-2} \), or (ii) a very small partial width \( \Gamma_{\sigma \to b\bar{b}}/m_{\sigma} \leq 10^{-4} \). Both of these cases are hardly conceivable.

Most probably, the mixing will be large enough for the unity in the denominator of Eqn. (11) to dominate. In that case, the observed Higgs branching ratio is \textit{independent} of the mixing angles, \( R_{h} \approx \frac{m_{h}^2}{m_{\sigma}^2} \times R_{\sigma} \). The physical reason for this is that both \( h \to gg \) and \( h \to b\bar{b} \) are dominated by the decay through mixing, \( h \to \sigma^{*} \to gg, b\bar{b} \).

### III. ESTIMATES OF \( \sigma \) DECAYS IN A POTENTIAL MODEL

We estimate the properties and decay widths of the \( \sigma \) particle in a potential model with a Yukawa potential. This particular framework is not necessarily realized in Nature but we remind ourselves that all we want is an existence proof for a situation where this \( \sigma \) might exist. The advantages of such a potential are: (1) it has two parameters, both of which can be reasonably closely related to parameters of the underlying theory (a new gauge coupling \( \alpha_{\text{new}} \) and a new symmetry breaking scale \( m \)), and (2) that the introduction of \( \mathbf{L} \cdot \mathbf{S} \) coupling allows for a low lying \( 2^{3}P_{0} \) state, while all other states become heavy or nonexistent.

We use such a potential model to calculate the toponium spectrum, to restrict the new coupling strength and screening mass, and to find the decay widths of the \( \sigma \). Especially, we use the usual potential model expression that relates the decay rate \( \sigma \to gg \) to the derivative of the radial wave function at the origin. To find these quantities we need to calculate the wave function of a particle moving in a Yukawa potential.
Consider the quantum mechanical problem of a point particle with mass $M$ moving in a Yukawa potential with screening mass $m$. The Schrödinger equation for the radial wave function is

$$\frac{u''_l(r)}{u_l(r)} = \frac{l(l + 1)}{r^2} - 2M \left( E + \frac{e^{-mr}}{r} \right). \quad (12)$$

In our case of a $t\bar{t}$ system $M$ is related by a rescaled reduced top mass by $2M = 4\pi\alpha_{new}m_t$. This relationship is at best only approximate, so that the parameter $\alpha_{new}$ above may not exactly coincide with the new gauge coupling. The binding energy of the $t\bar{t}$ bound state is then $E = 4\pi\alpha_{new}E$.

The energy levels in the Yukawa potential have been calculated with high precision in [8, 9]. Because in these references the wave functions are not quoted, we repeated the calculation of Ref. [8], where $u_l(r)$ is approximated by a linear combination of functions $r^{k+1}e^{-\beta r}$. Then the coefficients of the linear combination and also $\beta$ are varied in order to minimize the energy of the state. The results are shown in Table I.

As a consistency check, observe that the P-wave wave function at the origin is found to be extremely small. Because in our approximation this is not automatic, its smallness indicates that our calculation is highly accurate. Moreover, we checked that in the $m = 0$ case (pure Coulomb potential) we reproduce the exact results within the accuracy indicated by the valid digits kept, and also that the approximation to the wave function is better than 0.2% for all $r$. To the same accuracy, all the energy levels in Table I coincide with the ones given in Refs. [8, 9].

### Table I: The energy levels in a Yukawa potential.

| Screening parameter | $m/M = 0$ | $m/M = 0.1$ | $m/M = 0.2$ |
|---------------------|-----------|-------------|-------------|
| $E_{1S}$            | -0.4999   | -0.4071     | -0.3268     |
| $E_{2S}$            | -0.1249   | -0.04993    | -0.1211     |
| $E_{2P}$            | -0.1250   | -0.04653    | -0.004102   |
| $R_{2P}(0)$         | 0         | -0.000,0521 | 0.00125     |
| $R'_{2P}(0)$        | 0.2041    | 0.1802      | 0.1063      |
TABLE II: The binding energy levels and the $\sigma \rightarrow gg$ width in a Yukawa potential. Recall that the screening parameter, $m_M \equiv m_{\pi}^2 \alpha_{\text{new}} m_t$, arises from a new symmetry breaking scale (related in turn to the mass of the new gauge boson masses.)


dm/\alpha_{\text{new}} \quad 0 \quad 0.10 \quad 0.16 \quad 0.18 \quad 0.20 \quad 0.22 \\

$E_{\text{bind}}(\sigma \equiv 2P)/\alpha_{\text{new}}^2$ \quad -1.7 TeV \quad -640 GeV \quad -270 GeV \quad -160 GeV \quad -56 GeV \quad 0 \\

$E_{\text{bind}}(1S)/\alpha_{\text{new}}$ \quad -6.9 TeV \quad -5.6 TeV \quad -4.9 TeV \quad -4.7 TeV \quad -4.5 TeV \quad -4.3 TeV \\

$(\frac{\Gamma_{\sigma \rightarrow gg}}{m_\sigma})^{1/5}/(\frac{m_{\sigma} \alpha_{\text{new}}}{m_{\sigma}})$ \quad 3.52 \quad 3.17 \quad 2.58 \quad 2.25 \quad 1.94 \quad N/A

Inspection of Table II shows that increasing values of the screening parameter will lessen the binding energies. The critical values at which each of the states disappears are given in Ref. [9]:

$M_{1S} = 1.91, \quad M_{2S} = 0.31, \quad M_{2P} = 0.22.$

This will impose an upper bound on the screening mass.

B. Application to $t\bar{t}$

We note first of all that the Yukawa potential we are considering will not bind the $b\bar{b}$ system, even if the right handed $b$ quarks participated in the new interaction. The size of such a state in a pure Coulomb potential would be determined by the inverse bottom mass, the Bohr radius being $a \sim \frac{1}{4\pi \alpha_{\text{new}} m_b}$. The screening in the Yukawa potential ($m_{\text{new}} \gg \mathcal{O}(100 \, \text{GeV})$) completely shuts off the interaction at these distances. Another way to see this is that the rescaled reduced mass in $b\bar{b}$ would be $M = 4\pi \alpha_{\text{new}} \frac{m_b}{2}$. But the new interaction must be screened much above $m \gg 100 \, \text{GeV}$, and in order for a bound state to form the screening parameter must satisfy $\frac{m}{M} < 1.91$. These constraints can be satisfied only with $\alpha_{\text{new}} \gg 2$, an interaction too strong to be realistic.

Now we use the Yukawa potential to understand the toponium spectrum as well as the decay of $\sigma$ into two gluons. The latter is calculated using the usual QCD formula for quarkonium ($\chi_0 \rightarrow gg$) decay [10], $\Gamma_{\sigma \rightarrow gg} = 96\alpha_{\text{new}}^2 \frac{m_\sigma^4}{m_\sigma^2}$. The results are shown in Table II.

From Table II we quickly deduce the restrictions on the parameters when our scenario can occur. Due to the fifth power in the expression for the $\sigma \rightarrow gg$ width, $\alpha_{\text{new}}$ is restricted from
above. (Intuitively, the high power of $\alpha_{\text{new}}$ arises because large couplings pull together the wave function closer to the origin and provide a large $|R'(0)|^2$.) Essentially at $\alpha_{\text{new}} > 1$, the decay width suddenly becomes much larger than $m_\sigma$, so that the $\sigma$ resonance is not formed.

On the other hand, we must have $m \gg 100 \text{GeV}$ if we want to avoid a large contribution to top physics, which requires $\alpha_{\text{new}} > 0.5$ (See the second row in Table III). One conceivably realistic value is $\alpha_{\text{new}} = 0.7$ and $m = 140 \text{GeV}$, which would result in a binding energy of $-80 \text{GeV}$, i.e. $m_\sigma = 270 \text{GeV}$ and $\Gamma_{\sigma gg} = 300 \text{GeV}$.

There is one problem with the above situation. The P-wave state is necessarily in the second shell (corresponding to a $2^3P_0$), and in the realistic parameter range the first shell ($1S$) is very tightly bounded. Its mass is actually negative, which triggers condensation of the vacuum. This is reminiscent to a topcolor or top condensation scenario [11], which is outside the scope of the present paper, where we consider electroweak symmetry breaking through the usual Higgs mechanism.

We should note that this conclusion is brought about by the assumption that the new interaction does not generates a spin flip. That this assumption is unrealistic is shown by the simple analogy to nuclear forces, where there is a large spin-orbit interaction which even changes the ordering of the energy levels.

Fortunately, the inclusion of a simple $L \cdot S$ interaction quickly remedies the situation. Assume that the Yukawa potential includes such a term (we use the same screening mass in all terms for simplicity)

$$V(r) = -g^2 e^{-mr} (\cos \gamma - L \cdot S \sin \gamma).$$

Then all the $n \cdot L J$ states are $L \cdot S$ eigenstates, so that they feel a pure Yukawa potential (but for each state the strength of the coupling it feels, $\alpha_{\text{eff}}$, is different, as shown in Table III). Because the state that is most sensitive to the spin-orbit interaction happens to be the $\sigma$ (i.e. $3P_0$), we can shift all other states to smaller binding energies as compared to $\sigma$.

Replacing the parameter $\alpha_{\text{new}} \rightarrow \alpha_{\text{eff}}$, the numbers in Table III remain in force. The first

| State | $1S_0$ | $3P_1$ | $3P_2$ | $3P_1$ | $3P_0$ |
|-------|-------|-------|-------|-------|-------|
| $L \cdot S$ | 0 | 0 | 0 | 0 | 0 |
| $\alpha_{\text{eff}} / \alpha_{\text{new}}$ | $\cos \gamma$ | $\cos \gamma$ | $\cos \gamma - \sin \gamma$ | $\cos \gamma + \sin \gamma$ | $\cos \gamma + 2 \sin \gamma$ |

TABLE III: The binding energy levels and the $\sigma \rightarrow b \bar{b}$ width in a Yukawa potential.
shell with only S-wave states feels a different, weaker potential with \( \alpha_{eff}(1S) = \frac{\alpha_{new}}{1 + 2 \tan \gamma} \).

If the potential is dominated by the spin-orbit interaction, the binding energy of the 1S states can be made small. Because the screening parameter is also affected by the new \( \alpha_{eff} \) (while the screening mass \( m = 140 \text{GeV} \) remains unchanged), in the numerical example given above, for \(-1 < \cos \gamma < 0.2\), the only remaining state is the \( \sigma \equiv 2^3P_0 \).

IV. MODIFICATIONS TO THE HIGGS SIGNALS

We have shown an example when the mixing of the Standard Model Higgs boson with a new heavy scalar introduces significant changes in the couplings of the Higgs, most notably leading to a sizable, perhaps dominant, Higgs-gluon interaction. In the Higgs mass region \( 115 \text{GeV} < m_h < 180 \text{GeV} \) that we are considering, this changes the relative importance of the various production channels as well as the decay modes. This alters the strategy needed to discover the Higgs boson in a significant way, both at the hadron and lepton colliders.

In future leptonic colliders the environment is favorably clean for the detection of the Higgs. The “gold plated” mode \( e^+e^- \rightarrow Z^* \rightarrow Zh \) followed by the Higgs decay \( h \rightarrow b\bar{b} \) seems to be the easiest to see. The mixing discussed in this paper would increase the \( Z \) plus two gluon jets final state. A large \( BR(h \rightarrow gg) \sim 1 \) branching ratio arises if the new scalar \( \sigma \) dominantly decays into two gluons. This would diminish the \( Zb\bar{b} \) signal so that the Higgs might be missed altogether in searches for a standard Higgs. This is also important for the LEP2 data with the search restricted to \( h \rightarrow \text{two } b\)-jets.

At Tevatron Run 2, the most promising mode for the SM Higgs boson in the mass range \( 110 \text{GeV} \leq m_h \leq 130 \text{GeV} \) is \( q\bar{q} \rightarrow W^* \rightarrow Wh \rightarrow Wb\bar{b} \). However, in our scenario, the dominant decay mode of \( h \) may be \( h \rightarrow gg \), thus reducing the \( h \rightarrow b\bar{b} \) branching ratio. So the signal in the \( Wb\bar{b} \) mode may be greatly reduced, making it unobservable due to the large QCD background. On the other hand, the Higgs production in the mode \( gg \rightarrow h \) will be greatly enhanced in this scenario. Thus, the signal in the \( gg \rightarrow h \rightarrow WW^* \) channel will now be greatly enhanced. This will extend the lower mass range coverage below the usual case of \( m_h \sim 135 \text{GeV} \) via this mode. The mode \( gg \rightarrow h \rightarrow \tau^+\tau^- \) was a borderline case in the SM. With the enhanced \( gg \rightarrow h \) production rate of our scenario, the signal in this mode may be observable if the \( \tau \) detection efficiency and the missing \( E_T \) resolution can be improved in the upgraded CDF and D0 detectors.
In the SM, $\sigma(gg \to H) \sim 0.7\,pb$, while $B(h \to \gamma\gamma) \sim 2.2 \times 10^{-3}$ for $m_H = 120\,GeV$. In our scenario, the Higgs branching ratio to two photons will be increased, and so will be the $\sigma(gg \to H)$ production cross section. Thus, the two photon mode may be detectable at Tevatron Run 2. A similar enhancement will occur for the two photon mode at the LHC.

While in most of this paper we discussed the case of a light Higgs boson, it is worthwhile to point out that in the case of a heavy Higgs ($m_h \geq 135\,GeV$, when the dominant decay is $h \to WW, WW^*$) the QCD background is less severe on hadronic colliders. Then gluon fusion is the dominant process in the Standard Model, $gg \to h \to W^+W^-, WW^*$, providing an excellent signature. This signal will be greatly enhanced in the proposed scenario.

V. CONCLUSIONS

We have presented a scenario in which the coupling of the SM Higgs boson to two gluons can be greatly enhanced due to mixing with a scalar $t\bar{t}$ bound state (bound by new strong interactions). We have presented a potential model in which such a scenario may be realized. Such a scenario has major implications for the Higgs search at Tevatron Run 2 in the mass range of $110\,GeV \leq m_h \leq 180\,GeV$. Production and decay channels are both affected. The $Wb\bar{b}$ final state at Tevatron Run 2 will be suppressed, because the $h \to b\bar{b}$ branching ratio is suppressed. The $WW^*$ mode will be enhanced due to an increase in the $gg \to h$ production cross section, so that this mode will be effective even for a Higgs mass range lower than $135\,GeV$. Because Higgs production via $gg \to h$ mode is enhanced, the $\tau^+\tau^-$ and perhaps also the $\gamma\gamma$ final states may be detectable. Finally, we point out that LEP2 may have missed the Higgs boson because the dominant final states were $Zgg$ instead of $Zb\bar{b}$.

After the completion of this paper another work was published [12], which discusses a similar effect on the gluonic couplings of the Higgs boson. The underlying model in that paper is based on extra dimensions and is very different from our scenario.
Acknowledgments

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