Bounds on Bosonic Topcolor

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Abstract

We consider the phenomenology of models in which electroweak symmetry breaking is triggered by new strong dynamics affecting the third generation and is transmitted to the light fermions via a fundamental Higgs doublet. While similar in spirit to the old bosonic technicolor idea, such ‘bosonic top-color’ models are allowed by current phenomenological constraints, and may arise naturally in models with large extra dimensions. We study the parameter space of a minimal low-energy theory, including bounds from Higgs boson searches, precision electroweak parameters, and flavor changing neutral current processes. We show that the model can provide a contribution to $D^0\overline{D^0}$ mixing as large as the current experimental bound.

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I. INTRODUCTION

In spite of the quantitative success of the standard model, the mechanism of electroweak symmetry breaking remains unclear. Only a few years ago, bosonic technicolor models provided a relatively unconventional approach to solving this problem \cite{1-3}: electroweak symmetry was broken dynamically by a fermion condensate triggered by new strong forces, while a fundamental scalar field was responsible for transmitting these effects to the fermions through ordinary Yukawa couplings. These models did not require a conventional extended technicolor sector, and hence were freed from the associated flavor changing neutral current (FCNC) problems. Unfortunately, precision electroweak constraints rule out bosonic technicolor models at least in models where the strong dynamics is QCD-like and the S-parameter can be reliably estimated \cite{4}.

In this letter, we point out that a very similar scenario, bosonic topcolor, also provides a very simple low-energy effective theory, but one that is not in conflict with electroweak constraints. In this scenario, electroweak symmetry is partly broken by new strong dynamics that affects fields of the third generation, as in conventional topcolor scenarios \cite{5,6}, while a weakly-coupled scalar doublet transmits the symmetry breaking to the fermions via Yukawa couplings. Since this scenario involves both a fundamental ($H$) and a composite ($\Sigma$) Higgs field that both contribute to electroweak symmetry breaking, the usual problematic relation \cite{5} between the dynamical top quark mass and the electroweak symmetry breaking scale is not obtained. The result is a viable two Higgs doublet model of type III, which we will show survives the bounds from flavor changing neutral current processes and may provide interesting flavor-changing signals as well.

The possibility of topcolor models involving fundamental scalars has been considered in Refs. \cite{7,8}. In these papers, however, the fundamental Higgs field was strongly coupled, and the authors considered whether the fundamental field itself could trigger the formation of a $t\bar{t}$ condensate. Here we introduce $H$ as a weakly-coupled field and investigate the phenomenological consequences.
It is worth pointing out that a philosophical objection to the original bosonic technicolor scenarios, and the bosonic topcolor models described here, is that strong dynamics was originally intended to eliminate the need for a fundamental Higgs field altogether, as well as the associated problem with quadratic divergences. Recent theoretical developments relating to the possibility of low-scale quantum gravity [9] renders these objections hollow: The presence of a low string scale eliminates the conventional desert so that nonsupersymmetric low-energy theories with fundamental scalars are not unnatural. Moreover, in this setting there are new origins for the strong dynamics, namely the exchange of a nonperturbatively large number of gluon Kaluza-Kleinder excitations [10]. While we will not consider an explicit extra-dimensional embedding of the scenario described here, it seems that these considerations make the investigation of models with dynamical electroweak symmetry breaking and fundamental scalar fields well motivated.

In the next section we will present a simple realization of the bosonic topcolor idea following the Nambu-Jona-Lasinio approach [5]. Our first model is non-generic in the sense that we do not specify the most general set of higher-dimension operators that could appear in an arbitrary high-energy theory. However, it does provide a very convenient framework for parameterizing and exploring the basic phenomenological features of the scenario. After considering the phenomenological bounds, we will describe how to study the same type of scenario in a more general effective field theory approach. While we will not consider every phenomenological detail in this letter, we hope to obtain an accurate overall picture of the allowed parameter space. Finally, we will discuss flavor changing signals for the model, notably a potential contribution to $D^0\bar{D}^0$ mixing that can be as large as the current experimental bound. We then summarize our conclusions.

II. MINIMAL BOSONIC TOPCOLOR

Our high-energy theory is defined by

$$\mathcal{L} = \mathcal{L}_H + \mathcal{L}_{NJL},$$

(2.1)
where
\[
\mathcal{L}_H = D_\mu H^\dagger D^\mu H - m_H^2 H^\dagger H - \lambda (H^\dagger H)^2 - h_t (\bar{\psi}_L H t_R + h.c.)
\]  
(2.2)

and
\[
\mathcal{L}_{NJL} = \frac{\kappa}{\Lambda^2} \bar{\psi}_L t_R t_R \psi_L
\]  
(2.3)

The field \( H \) is a fundamental scalar doublet, and \( \Lambda \) characterizes the scale at which new physics is present that generates the nonrenormalizable interaction in Eq. (2.3). In light of our introductory remarks, we will assume henceforth that \( \Lambda \lesssim 100 \text{ TeV} \). In this minimal scenario we assume that the right-handed top, and left-handed top-bottom doublet \( \psi_L \) experience the new strong interactions. Immediately beneath the scale \( \Lambda \) we may rewrite Eqs. (2.2) and (2.3) as
\[
\mathcal{L} = D_\mu H^\dagger D^\mu H - m_H^2 H^\dagger H - \lambda (H^\dagger H)^2 - c \Lambda^2 \Sigma^\dagger \Sigma
\]
\[-h_t (\bar{\Psi}_L t_R H + h.c.) - g_t (\bar{\Psi}_L t_R \Sigma + h.c.)
\]  
(2.4)

where \( \Sigma \) is a non-propagating auxiliary field. Using the equations of motion, \( \Sigma = -g_t (\bar{t}_R \psi_L)/(c \Lambda^2) \) and one recovers Eqs. (2.2) and (2.3) with the identification \( \kappa = g_t^2/c \).

At energies \( \mu \ll \Lambda \), quantum corrections induce a kinetic term for \( \Sigma \), so that it becomes a dynamical field, a composite Higgs doublet in the low-energy theory. In order to study the quantum corrections to Eq. (2.4) it is convenient for us to define the column vector
\[
\Phi = \begin{pmatrix} \Sigma \\ H \end{pmatrix}
\]  
(2.5)

Then the kinetic term at the scale \( \mu \) may be written \( \partial_\mu \Phi^\dagger Z \partial^\mu \Phi \), with
\[
Z = \begin{pmatrix}
\frac{g_t^2 N_C \ln(\Lambda/\mu)}{8\pi^2} & \frac{g_t h_t N_C \ln(\Lambda/\mu)}{8\pi^2} \\
\frac{g_t h_t N_C \ln(\Lambda/\mu)}{8\pi^2} & 1 + \frac{h_t^2 N_C \ln(\Lambda/\mu)}{8\pi^2}
\end{pmatrix}
\approx \begin{pmatrix}
1 & 0 \\
0 & 1
\end{pmatrix}
\]  
(2.6)

In general, \( Z \) must be diagonalized and rescaled so that the kinetic terms assume the canonical form. However, in most of the parameter space that we consider later in this paper
$h_t$ is small enough that the off-diagonal elements of $Z$ are numerically irrelevant; thus we use the simpler approximate form parameterized by $r$ in Eq. (2.6). Properly normalized kinetic terms are then obtained via the substitution $\Sigma \to r\Sigma$. Quantum corrections also induce quartic self interactions, and mixing in the $\Phi$ mass matrix. We retain the largest self-coupling, $(\Sigma^\dagger\Sigma)^2$ with coefficient $\lambda_\Sigma = g_t^4 N_C \ln(\Lambda/\mu)/(4\pi^2)$; the $\Phi$ mass matrix is given by

$$L_{mass} = -\Phi^\dagger \mathcal{M}^2 \Phi,$$  

(2.7)

with

$$\mathcal{M}^2 = \begin{pmatrix} r^2 m_\Sigma^2 & r \delta m^2 \\ r \delta m^2 & m_H^2 \end{pmatrix},$$  

(2.8)

where $\delta m^2 = -\frac{N_c}{8\pi^2} g_t h_t \Lambda^2$. Eq. (2.8) reflects the fact that both the diagonal and off-diagonal entries receive quadratically divergent radiative corrections. For the diagonal elements, the tree-level mass terms present in Eq. (2.4) can be fine-tuned against these radiative corrections (as in the standard model) so that $m_\Sigma$ and $m_H$ are well beneath the cutoff scale $\Lambda$. On the other hand, there is no tree-level $H\Sigma$ mass mixing term given the way we defined our high-energy theory in Eqs. (2.2)-(2.3). However, since we are considering the situation where the scale $\Lambda$ is relatively low ($< 100$ TeV) and where the coupling $h_t$ is small (see Figs. 1 and 2), the off-diagonal elements will also be much smaller than the cutoff. For the case where such tree-level mass mixing is present at the scale $\Lambda$, the reader should refer to Section 3.

Electroweak symmetry will be broken in this model if $\Sigma$ and/or $H$ acquire vacuum expectation values (vevs). There are several ways this can happen depending on the values of the different parameters in the model. We are principally interested in the case where $m_H^2 > 0$, so that electroweak symmetry breaking is triggered by the strong dynamics and the vev of $H$ can be interpreted as a subsidiary effect. Thus, it is necessary to study the scalar potential,

$$V(\Sigma, H) = r^2 m_\Sigma^2 \Sigma^\dagger \Sigma + m_H^2 H^\dagger H + r \delta m^2 \left( \Sigma^\dagger H + h.c. \right) + \lambda \left( H^\dagger H \right)^2 + \lambda_\Sigma r^4 \left( \Sigma^\dagger \Sigma \right)^2.$$  

(2.9)
Rather than search directly for minima in a five-dimensional parameter space \((m^2_\Sigma, m^2_H, \delta m^2, \lambda, \lambda_\Sigma)\) we extremize the potential and solve for \(m^2_\Sigma\) and \(m^2_H\) in terms of the \(\Sigma\) and \(H\) vevs. It is much more manageable to study the remaining constrained three-dimensional parameter space and determine which points correspond to stable local minima. If we denote the vevs of \(\Sigma\) and \(H\) by \(v_1/\sqrt{2}\) and \(v_2/\sqrt{2}\), one finds

\[
m^2_\Sigma = -\frac{1}{rv_1} \left( \delta m^2 v_2 + \lambda_\Sigma r^3 v_1^3 \right),
\]

(2.10)

\[
m^2_H = -\frac{1}{v_2} \left( \delta m^2 rv_1 + \lambda v_2^3 \right).
\]

(2.11)

From Eq. (2.9), one may obtain the mass matrices for the scalars, pseudoscalars, and charged scalars:

\[
M_S = \frac{1}{2} \begin{pmatrix}
m^2_\Sigma r^2 + 3\lambda_\Sigma r^4 v_1^2 & \delta m^2 r \\
\delta m^2 r & m^2_H + 3\lambda v_2^2
\end{pmatrix},
\]

(2.12)

\[
M_P = \frac{1}{2} \begin{pmatrix}
m^2_\Sigma r^2 + \lambda_\Sigma r^4 v_1^2 & \delta m^2 r \\
\delta m^2 r & m^2_H + \lambda v_2^2
\end{pmatrix},
\]

(2.13)

\[
M_+ = \begin{pmatrix}
m^2_\Sigma r^2 + \lambda_\Sigma r^4 v_1^2 & \delta m^2 r \\
\delta m^2 r & m^2_H + \lambda v_2^2
\end{pmatrix}.
\]

(2.14)

The Higgs field vevs are responsible for producing the proper gauge boson masses, \textit{i.e.}

\[
v_1^2 + v_2^2 = (246 \text{ GeV})^2,
\]

(2.15)

as well as the mass of the top quark

\[
m_t = (r g_t v_1 + h_t v_2)/\sqrt{2}.
\]

(2.16)

This expression shows that the top quark receives both an ordinary and a dynamical contribution. Since we focus on small values of \(h_t\) in this letter, the top quark mass is mostly dynamical, originating from the first term in Eq. (2.16). In this limit, the vevs \(v_1\) and \(v_2\) are determined by the choice of scales \(\Lambda\) and \(\mu\), since the quantity \(r g_t\) is independent of \(g_t\).
III. PHENOMENOLOGY

Notice that all the freedom in Eqs. (2.10)-(2.14) is fixed by specifying $\Lambda$, $\mu$, $h_t$ and $\lambda$. Thus, for a fixed choice of $\Lambda < 100$ TeV and $\mu$ of order the weak scale, we may map our results onto the $\lambda$-$h_t$ plane. Fig. 1 displays results for $\Lambda = 10$ TeV and Fig. 2 for $\Lambda = 100$ TeV, with $g_t = 1$. In each case, the intersecting solid lines indicate where $m_\Sigma^2$ or $m_H^2$ change sign, with positive values lying above the corresponding line. Figs. 1a and 2a provide mass contours for the lightest neutral scalar and charged scalar states; Figs. 1b and 2b display constant contours for the electroweak parameters $S$ and $T$. These were computed using formulae available in the literature for general two Higgs doublet models [11],

\[ S = \frac{1}{12\pi} \left( s_{\alpha-\beta}^2 \ln \frac{M_2^2}{M_H^2} + g(M_1^2, M_2^2) - \frac{1}{2} \ln \frac{M_1^2}{M_2^2} - \frac{1}{2} \ln \frac{M_2^2}{M_3^2} \right) \]

\[ + c_{\alpha-\beta}^2 \left( \ln \frac{M_1^2}{M_H^2} + g(M_2^2, M_3^2) - \frac{1}{2} \ln \frac{M_2^2}{M_3^2} - \frac{1}{2} \ln \frac{M_3^2}{M_4^2} \right), \quad (3.1) \]

and

\[ T = \frac{3}{48\pi s_w^2 m_W^2} \left( s_{\alpha-\beta}^2 \left[ f(M_1^2, M_2^2) + f(M_3^2, M_4^2) - f(M_1^2, M_3^2) \right] \right) \]

\[ + c_{\alpha-\beta}^2 \left[ f(M_2^2, M_4^2) + f(M_3^2, M_4^2) - f(M_2^2, M_3^2) \right], \quad (3.2) \]

where $M_1$, $M_2$, $M_3$, and $M_4$ are the light scalar, heavy scalar, pseudoscalar, and charged scalar masses respectively, and $\beta = \tan^{-1}(v_2/v_1)$. The scalar mixing angle $\alpha$ and the functions $f$ and $g$ are defined in Ref. [11]. Figs. 1c and 2c show regions excluded by (i) the current LEP bound on neutral Higgs production, (ii) bounds on the $S$ and $T$ parameters, (iii) bounds on the charged scalar mass from $b \to s\gamma$. In the first case, we compute the production cross section for the lightest scalar state $\phi_s$,

\[ \sigma(e^+e^- \to Z\phi_s) = s_{\alpha-\beta}^2 \sigma_{SM}(e^+e^- \to ZH^0) \]

(3.3)

and compare to the corresponding standard model cross section for a Higgs boson with mass equal to the current LEP bound, $m_H < 107.9$ GeV, 95% CL [12]. In the case of the $S$ and $T$ bounds, we consider the results of global electroweak fits quoted in the Review...
of Particle Properties [13], \( S = -0.26 \pm 0.14 \) and \( T = -0.11 \pm 0.16 \) [4], which assume a reference Higgs mass of 300 GeV. We show the two standard deviation limit contours for \( S \) and \( T \) separately wherever an upper or lower limit is exceeded. (Note that we don’t take into account correlations between \( S \) and \( T \) in determining this exclusion region.) Finally, we plot the charged Higgs mass limit \( m_{H^+} > \left[ 244 + 63/(\tan \beta)^{1.3} \right] \) GeV from \( b \to s\gamma \) [14]. This is the strongest, albeit indirect, charged Higgs mass limit listed in Ref. [13]. Although, strictly speaking, this bound applies to a type II two-Higgs doublet model, the leading top quark loop contribution is the same in our model; the top quark-charged scalar coupling is given by

\[
\Phi^+ \frac{g}{2\sqrt{2}M_W} \left[ 7(m_t^H \cot \beta - m_t^\Sigma \tan \beta)V_{tq}(1 - \gamma^5)q - 7\cot \beta V_{tq}m_q^H(1 + \gamma^5)q \right]
\]

in the case where the Cabibbo-Kobayashi-Maskawa (CKM) matrix \( V \) originates from diagonalization of the down quark Yukawa matrix alone (the reason for this assumption is given in the following section). Here \( m_t^H \) and \( m_t^\Sigma \) refer to contributions to the top mass from the \( H \) and \( \Sigma \) vevs, respectively. For most of the parameter range of interest to us, \( m_t^H \ll m_t^\Sigma \) and the interaction in Eq. (3.4) reduces to that of a type II model, and the \( b \to s\gamma \) bound is approximately valid. In both Figs. 1 and 2 a rectangular region is shown in which the charged scalars are heavy enough to weaken the flavor changing neutral current bounds, without exceeding that of the \( T \) parameter.

**IV. FLAVOR CHANGING SIGNALS**

The fact that one of our two Higgs doublets (\( \Sigma \)) couples preferentially to the top quark leads to a potentially interesting source of flavor violation in the model. While the charge \(-1/3\) quark masses and neutral scalar interactions both originate via couplings to \( H \) (and hence are simultaneously diagonalizable), the same is not true in the charge \( 2/3 \) sector, where the mass matrix depends on both the \( H \) and \( \Sigma \) vevs,
\begin{equation}
M^U = Y^U v_1 \frac{v_1}{\sqrt{2}} + Y^U v_2 \frac{v_2}{\sqrt{2}}.
\end{equation}

For concreteness, let us consider a definite Yukawa texture:

\begin{equation}
M^U = \begin{pmatrix}
0 & 0 & 0 \\
0 & 0 & 0 \\
0 & r g_t & 0
\end{pmatrix} \frac{v_1}{\sqrt{2}} + \begin{pmatrix}
\lambda^8 & \lambda^5 & \lambda^3 \\
\lambda^5 & \lambda^4 & \lambda^2 \\
\lambda^3 & \lambda^2 & h_t
\end{pmatrix} \frac{v_2}{\sqrt{2}}.
\end{equation}

Here \( \lambda = 0.22 \) is the Cabibbo angle, and we have picked a symmetric texture for the fundamental Higgs Yukawa couplings that approximately reproduces the correct CKM angles. Dropping the factors of \( v_i/\sqrt{2} \), the matrices shown give the neutral scalar couplings in our original field basis. In the quark (and Higgs) mass eigenstate basis, there will be flavor-changing top and charm quark interactions. Here, we focus only on the latter. CKM-like rotations that diagonalize the mass matrices yield 1-2 neutral scalar couplings of order \( \lambda^5 \).

It is straightforward to estimate the contribution to \( D^0-\bar{D}^0 \) mixing,

\begin{equation}
|\frac{\Delta m_D}{m_D}|_{\text{new}} \approx \lambda^{10} \frac{f_D^2}{12 M_{\phi}^2} \left[-1 + 11 \frac{m_D^2}{(m_c + m_u)^2}\right].
\end{equation}

For \( f_D \approx 200 \text{ MeV} \), this contribution saturates the current experimental bound, \( \Delta m_D < 1.58 \times 10^{-10} \text{ MeV} \) \cite{13}, for \( M_{\phi} \lesssim 495 \text{ GeV} \). The reason that we do not include this as a bound is that the 1-2 neutral scalar couplings need not be \( O(\lambda^5) \); they could in principle be zero, if the CKM matrix results from the diagonalization of the down quark Yukawa matrix alone. Since this is the least constrained possibility, we adopted this assumption in Eq. (3.4) for computing the \( b \to s \gamma \) exclusion region. Generically, however, we see that bosonic topcolor models predict significant contributions to \( D^0-\bar{D}^0 \) mixing, potentially as large as the current experimental bound.

\section*{V. GENERALIZATIONS}

The scenario described in the previous section is particularly convenient in that the basic phenomenology can be described in a two-dimensional parameter space, for fixed \( \Lambda \)
and $\mu$. However, a realistic high-energy theory is likely to provide more than the single higher-dimension operator in Eq. (2.3). In this section we briefly describe the effective field theory approach for constructing the most general low-energy effective bosonic topcolor theory. Given our assumption that $\psi_L$ and $t_R$ experience the new strong dynamics, the strongly-coupled sector of the theory possesses a global symmetry $G = SU(2)_L \times U(1)_R$, that is spontaneously broken by the $t\bar{t}$ condensate to the $U(1)$ that counts top quark number. If we denote the elements of this $SU(2) \times U(1)$ by $U$ and $V$, respectively, then the transformation properties of the fields are given by

$$\psi_L \rightarrow U\psi_L, \quad t_R \rightarrow Vt_R, \quad \text{and} \quad \Sigma \rightarrow U\Sigma V^\dagger, \quad (5.1)$$

where $V$ is a phase. The Yukawa couplings of the fundamental Higgs field explicitly break $G$, so we may treat $h_t H$ as a ‘spurion’ transforming as

$$h_t H \rightarrow U(h_t H)V^\dagger. \quad (5.2)$$

We may now include $h_t H$ systematically in an effective Lagrangian by forming all possible $G$-invariant terms. At the renormalizable level,

$$L_{\text{eff}} = D^\mu H^\dagger D_\mu H + D^\mu \Sigma^\dagger D_\mu \Sigma$$

$$- m_H^2 H^\dagger H - m_\Sigma^2 \Sigma^\dagger \Sigma + m_{H\Sigma} h_t (H^\dagger \Sigma + h.c.)$$

$$- \lambda (H^\dagger H)^2 - \lambda_0 (\Sigma^\dagger \Sigma)^2 + h_t (h^\dagger \Sigma)\Sigma^\dagger \Sigma + \cdots$$

$$- h_t \overline{\psi}_L H t_R - g_{\Sigma} \overline{\psi}_L \Sigma t_R + h.c. \quad (5.3)$$

Note that we have eliminated a possible kinetic mixing term by field redefinitions, which do not affect the form of the other terms. The $\cdots$ represent all the other possible quartic terms which are higher order in $h_t$. Unlike the model described in the previous section, we no longer have a boundary condition at the scale $\Lambda$ that sets $\lambda_0(\Lambda) = 0$ and $m_{H\Sigma}(\Lambda) = 0$, thus introducing two additional degrees of freedom into the scalar potential. Since we are now working directly with the low-energy theory, the scale $\Lambda$ is not input directly, but rather can be computed by determining the scale at which the wavefunction renormalization of the
Σ field vanishes. At this scale, Σ again becomes an auxiliary field, and may be eliminated using the equations of motion, leaving a more general set of higher-dimension operators than we had assumed originally in Eq. \(2.3\).

A complete investigation of the parameter space of this generalized model is beyond the scope of this letter. Before closing, we point out that there are reasonable parameter choices in Eq. \(5.3\) where the resulting phenomenology is similar to the minimal model considered in Section 2. In Fig. 3 we provide the same information given in Figs. 1 and 2 for the general bosonic topcolor model, with \(m_{H^2}^2 = (400 \text{ GeV})^2\) and \(\lambda_0 = 1\). It is interesting that in this case the allowed band delimited by the FCNC and \(T\) parameter lines lies mostly in the region where both \(m_{H^2}^2\) and \(m_{H^2}^2\) are positive; in this region the mixing term in Eq. \(2.8\) drives one of the scalar mass squared eigenvalues negative so that electroweak symmetry is broken. A full exploration of this parameter space will be provided elsewhere \[15\].

VI. CONCLUSIONS

In this letter, we have described models in which electroweak symmetry breaking is triggered by strong dynamics affecting the third generation but transmitted to the fermions by a weakly-coupled, fundamental Higgs doublet. We have argued in the Introduction that such models are not unnatural given recent developments in low-scale quantum gravity. Our minimal scenario, while probably not representing the ultimate high-energy theory, has the virtue of allowing a simple parameterization of the basic phenomenology of the model. It is our hope that others will adopt it as the basis for further phenomenological study. Issues that one could address include relaxation of our small \(h_t\) approximation, flavor-changing top quark processes, and the collider physics of the model. We also described how the scenario may be generalized using effective field theory techniques. Unlike bosonic technicolor models, bosonic topcolor is not excluded by current phenomenological bounds. Moreover, the model has interesting flavor-changing signals such as a contribution to \(D^0-\bar{D^0}\) mixing that could be as large as current experimental bounds.
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FIG. 1. Minimal model, $\Lambda = 10$ TeV. (a) Neutral (dashed) and charged (dotted) mass contours, units of GeV, (b) $S$ (dotted) and $T$ (dashed) parameter contours, (c) Exclusion regions.
FIG. 2. Same as Fig. 1, with $\Lambda = 100$ TeV.
FIG. 3. General model, $m_H^2 = (400 \text{ GeV})^2$, $\lambda_0 = 1$. Notation is the same as in Figs. 1 and 2.