

Hidden extra $U(1)$ at the electroweak/TeV scaleB. N. Grossmann,^{1,*} B. McElrath,^{2,†} S. Nandi,^{1,3,‡} and Santosh Kumar Rai^{1,§}¹*Department of Physics and Oklahoma Center for High Energy Physics, Oklahoma State University, Stillwater, Oklahoma 74078, USA*²*Universität Heidelberg, 16 Philosophenweg, 69120 Heidelberg, Germany*³*Fermi National Accelerator Laboratory, P.O. Box 500, Batavia, Illinois 60510, USA^{||}*

(Received 19 July 2010; published 24 September 2010)

We propose a simple extension of the standard model (SM) by adding an extra $U(1)$ symmetry which is hidden from the SM sector. Such a hidden $U(1)$ has not been considered before, and its existence at the TeV scale can be explored at the LHC. This hidden $U(1)$ does not couple directly to the SM particles, and couples only to new $SU(2)_L$ singlet exotic quarks and singlet Higgs bosons, and is broken at the TeV scale. The dominant signals at the high-energy hadron colliders are multilepton and multi- b -jet final states with or without missing energy. We calculate the signal rates as well as the corresponding standard model background for these final states. A very distinctive signal is 6 high p_T b -jets in the final state with no missing energy. For a wide range of the exotic quarks masses the signals are observable above the background at the LHC.

DOI: 10.1103/PhysRevD.82.055021

PACS numbers: 12.60.Cn, 12.60.Fr, 12.90.+b

I. INTRODUCTION

The standard model (SM) is now well established as the effective theory of the quarks, leptons, and the gauge bosons below the TeV scale. However, it is almost universally believed that the SM is not the final theory. Many extensions of the SM have been proposed to solve some of the shortcomings of the SM. Grand unified theories (GUT's) have been proposed to unify the three gauge couplings into one. Supersymmetric extensions have been proposed to solve the gauge hierarchy problem. A singlet Higgs boson, with a Z_2 symmetry, has been added to the SM which can serve as a plausible candidate for dark matter [1]. One or more extra spacelike compact dimensions has been added to the usual four dimensions to incorporate TeV scale as the fundamental scale of gravity [2], or to unify the gauge and Higgs bosons, as well as the fermions, and to understand the Yukawa interactions as part of the gauge interactions [3]. Most of these extensions involve new gauge interactions, commonly an extra $U(1)$, as well as new particles beyond the SM.

Many kinds of extra $U(1)$ gauge symmetries have been considered. These include the left-right symmetric model [4], $SO(10)$ or E_6 GUT's, superstring E_6 models [5], top-flavor models [6], and string-inspired supersymmetric models. [7] In most of these models, the SM fermions and the electroweak (EW) Higgs boson carry nontrivial charges under the $U(1)$. Other variations of the extra $U(1)$ symmetry, such as a hadrophobic $U(1)$, leptophobic $U(1)$, and an extra $U(1)$ which couples only to the third family of fermions, have been considered [5].

Hidden sectors of matter are ubiquitous among models due to the need to break supersymmetry, as well as the common addition of particle dark matter which cannot be charged under the standard model gauge groups unless it is very heavy [8]. An extra $U(1)$ can be a natural way to link the “dark” sector with the standard model.

In this work, we consider an extra $U(1)$ symmetry [9] in which the SM particles (the SM fermions, gauge bosons, and the EW doublet Higgs bosons) are neutral. We call this a hidden extra $U(1)$. Only new exotic quarks and the EW singlet Higgs bosons couple to this extra $U(1)$ gauge boson. These exotic quarks and the EW singlet Higgs bosons act as messenger particles between the hidden $U(1)$ sector and the SM sector. This extra $U(1)$ symmetry is broken at the EW scale by the vacuum expectation value (VEV) of the EW singlet Higgs boson. Thus this extra gauge boson, the exotic quarks, and the singlet Higgs boson all acquire masses at the EW scale, and can be searched for at high-energy colliders, such as the Tevatron and the LHC. The dominant signals of our scenario at the hadron colliders are multi- b and multilepton final states, with little or no missing energy. The paper is organized as follows. In Sec. II we present the model with the extra $U(1)$ hidden symmetry and discuss the new particle content of the model and their interactions with the SM particles arising out of mixings. In Sec. III we summarize the phenomenological implications of the model and study the experimental signatures, where we present our analysis and the first study of a six- b final state and its SM background at LHC. We finally summarize and conclude in Sec. IV.

II. THE HIDDEN $U(1)$ MODEL

Our gauge symmetry is the usual standard model $SU(3)_C \times SU(2)_L \times U(1)_Y$ supplemented by an extra

*benjamin.grossmann@okstate.edu

†bob.mcelrath@cern.ch

‡s.nandi@okstate.edu

§santosh.raai@okstate.edu

^{||}S. Nandi was a Summer Visitor at Fermilab

$U(1)$ symmetry, which we call $U(1)'$. We introduce two exotic quarks D_L and D_R which are color triplets and singlets under the $SU(2)_L$ gauge symmetry, but charged under the $U(1)_{Y'}$. We denote the gauge boson for the $U(1)'$ by Z' . We also introduce a complex Higgs field S which is a color and EW singlet, and has a charge q' under the $U(1)'$. This singlet Higgs field has a VEV v_S at the TeV or EW scale, and breaks the $U(1)'$ symmetry.

The Lagrangian for the gauge part of the interaction for the exotic D quark is given by the usual gauge interaction under the $SU(3)_C$ symmetry with the gauge coupling g_3 . The EW and $U(1)'$ interactions of the matter fields with the gauge bosons can be obtained from the Lagrangian

$$\mathcal{L} = \bar{q}_L^i i \mathcal{D}_2 q_L^i + \bar{\ell}_L^i i \mathcal{D}_2 \ell_L^i + \bar{u}_R^i i \mathcal{D}_1 u_R^i + \bar{d}_R^i i \mathcal{D}_1 d_R^i + \bar{e}_R^i i \mathcal{D}_1 e_R^i + \bar{D} i \mathcal{D}_1 D, \quad (1)$$

where q_L, ℓ_L are the $SU(2)_L$ quark and lepton doublets while u_R, d_R, e_R , and D are the $SU(2)_L$ up-quark, down-quark, lepton, and exotic quark singlets, respectively. The different covariant derivatives are defined as

$$\begin{aligned} \mathcal{D}_{2\mu} &= \partial_\mu - i \frac{g_2}{2} \tau \cdot W_\mu - i \frac{g'}{2} Y B_\mu, \\ \mathcal{D}_{1\mu} &= \partial_\mu - i \frac{g'}{2} Y B_\mu, \\ \mathcal{D}'_{1\mu} &= \partial_\mu - i \frac{g'}{2} Y B_\mu - i g_{Z'} Y_{Z'} Z'_\mu, \end{aligned} \quad (2)$$

where τ 's are the Pauli matrices; $Y, Y_{Z'}$ are the charges of the matter fields under the $U(1)_Y$ and the new gauge group $U(1)'$ respectively; while Z' represents the new gauge boson.

The Higgs potential, with the usual doublet Higgs H and the EW singlet Higgs S , is

$$V(H, S) = -\mu_H^2 (H^\dagger H) - \mu_S^2 (S^\dagger S) + \lambda_H (H^\dagger H)^2 + \lambda_{HS} (H^\dagger H) (S^\dagger S) + \lambda_S (S^\dagger S)^2. \quad (3)$$

We can also write a mass term for the vectorlike quark,

$$\mathcal{L}_{\text{mass}} = M_D \bar{D}_L D_R. \quad (4)$$

Note that the new exotic vectorlike quark D is like a new flavor, and it has color, hypercharge, and an extra $U(1)'$ interaction, but no $SU(2)_L$ interaction. Since this new D quark is vectorlike with respect to both $U(1)$ as well as $U(1)'$, the model is anomaly free. Without any other interaction, the D quark will be stable. As none of the SM particles are charged under the new $U(1)'$ symmetry, the new symmetry will remain hidden from the SM, provided the gauge-kinetic-mixing terms are strongly suppressed. However, its gauge quantum numbers allow flavor changing Yukawa interactions with the bottom, strange, and down quarks via the singlet Higgs boson S

$$\begin{aligned} \mathcal{L}_{\text{Extra Yukawa}} &= Y_{Db} \bar{D}_L b_R S + Y_{Ds} \bar{D}_L s_R S \\ &+ Y_{Dd} \bar{D}_L d_R S + \text{H.c.} \end{aligned} \quad (5)$$

Note that in order for the above Lagrangian to be hypercharge singlet, the hypercharge of both D_L and D_R must be equal to that of b_R . This also requires that the $U(1)'$ charge ($Y_{Z'}$) for the exotic quark D must satisfy $Y_{Z'} = q'$. Such a term in the Lagrangian leads to mixing between the down-type quarks with the new exotic vectorlike quark D , giving rise to EW decay modes for the heavy quark.

We assume that the parameters in the Higgs potential are such that H has VEV at the EW and S has a VEV around the TeV scale. Then, in the unitary gauge, we can write the H and S fields as

$$H(x) = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_H + H_0(x) \end{pmatrix}, \quad S(x) = \frac{1}{\sqrt{2}} (v_S + S_0(x)), \quad (6)$$

where v_H is the VEV of the doublet Higgs and v_S is the singlet VEV. From the minimization of the Higgs potential, we obtain

$$v_H^2 = \frac{\mu_H^2 - \frac{\lambda_{HS}}{2\lambda_S} \mu_S^2}{\lambda_H - \frac{\lambda_{HS}^2}{4\lambda_S}}, \quad v_S^2 = \frac{\mu_S^2 - \frac{\lambda_{HS}}{2\lambda_H} \mu_H^2}{\lambda_S - \frac{\lambda_{HS}^2}{4\lambda_H}}. \quad (7)$$

The scalar mass-squared matrix in the (H_0, S_0) basis is given by

$$\mathcal{M}^2 = \begin{pmatrix} 2\lambda_H v_H^2 & \lambda_{HS} v_H v_S \\ \lambda_{HS} v_H v_S & 2\lambda_S v_S^2 \end{pmatrix}. \quad (8)$$

The masses of the two mass eigenstate Higgs scalars ϕ_H and ϕ_S as well as their mixing angle β , in terms of the fundamental parameters of the Lagrangian, can be obtained from the above mass matrix. In particular, the mixing angle β is given by

$$\tan 2\beta = \frac{\lambda_{HS} v_H v_S}{\lambda_S v_S^2 - \lambda_H v_H^2}. \quad (9)$$

In addition to the usual gauge interaction for the H_0 and S_0 , the interaction among the Higgs from the potential $V(H_0, S_0)$ after symmetry breaking is given by

$$\begin{aligned} V(H_0, S_0) &= \lambda_H v_H H_0^3 + \frac{\lambda_{HS} v_S}{2} (H_0^2 S_0 + H_0 S_0^2) \\ &+ \lambda_S v_S S_0^3 + \frac{\lambda_H}{4} H_0^4 + \frac{\lambda_{HS}}{4} H_0^2 S_0^2 \\ &+ \frac{\lambda_S}{4} S_0^4. \end{aligned} \quad (10)$$

The interaction among the Higgs mass eigenstates ϕ_H and ϕ_S can be obtained by using

$$\begin{aligned} H_0 &= \phi_H \cos \beta + \phi_S \sin \beta, \\ S_0 &= -\phi_H \sin \beta + \phi_S \cos \beta. \end{aligned} \quad (11)$$

To explore the phenomenological implications of the model, we need to consider the various mixings which lead to the effective interaction of these exotic particles to SM particles and are responsible for their decays. We have already considered the mixing in the scalar sector

of the model which has interesting consequences for Higgs searches at colliders such as LHC [10]. We also find that by allowing the Yukawa interactions given in Eq. (5), there will be mixing between the down-type quarks with the new exotic quark D once the scalar S gets a VEV. The mixing between the down-type quarks with the exotic D quark gives rise to EW decay modes for the heavy quark. The heavy Z' also has additional interactions which lead to interesting decay modes.

For simplicity, we assume that only Y_{Db} is nonzero in Eq. (5) while the other Yukawa coefficients are negligibly small. This would imply that the exotic quark mixes only with the bottom quark, thus indirectly affecting the top-bottom vertex (V_{tb} in the Cabibbo-Kobayashi-Maskawa matrix) as well as inducing a coupling between the exotic D quark with the top quark. The mixing can be parametrized in terms of two mixing angles θ_L and θ_R , which represent the mixing angles of the b_L and b_R with D_L and D_R , respectively. Expressing the gauge eigenstates for the mixing quarks as b^0 and D^0 , the mass matrix in the (b^0, D^0) basis is given by

$$\mathcal{M} = \begin{pmatrix} y_b v_H / \sqrt{2} & 0 \\ Y_{Db} v_S / \sqrt{2} & M_D \end{pmatrix}. \quad (12)$$

This matrix can be diagonalized with a biunitary transformation $\mathcal{M}_{\text{diag}} = \mathcal{R}_L \mathcal{M} \mathcal{R}_R^\dagger$, where \mathcal{R}_L and \mathcal{R}_R are unitary matrices which rotate the left-chiral and right-chiral gauge eigenstates to the mass eigenstates, respectively. The interaction of the physical mass eigenstates (b, D) can then be obtained by writing the gauge basis states as

$$\begin{aligned} b_i^0 &= b_i \cos \theta_i + D_i \sin \theta_i, \\ D_i^0 &= -b_i \sin \theta_i + D_i \cos \theta_i. \end{aligned} \quad (13)$$

The rotation matrices \mathcal{R}_i are given by

$$\mathcal{R}_i = \begin{pmatrix} \cos \theta_i & \sin \theta_i \\ -\sin \theta_i & \cos \theta_i \end{pmatrix}, \quad \text{where } i = L, R. \quad (14)$$

The corresponding mixing angles for the left- and right-handed fields follow from diagonalizing the matrices $\mathcal{M} \mathcal{M}^\dagger$ and $\mathcal{M}^\dagger \mathcal{M}$, respectively, and are given by

$$\begin{aligned} \tan 2\theta_L &= \frac{-2Y_{Db} y_b v_S v_H}{2M_D^2 + Y_{Db}^2 v_S^2 - y_b^2 v_H^2}, \\ \tan 2\theta_R &= \frac{-2\sqrt{2}Y_{Db} v_S M_D}{2M_D^2 - Y_{Db}^2 v_S^2 - y_b^2 v_H^2}. \end{aligned} \quad (15)$$

III. PHENOMENOLOGICAL IMPLICATIONS

In hadronic colliders such as the LHC and Tevatron, the dominant signals arise from the pair productions of the exotic colored quarks, D and \bar{D} , and their subsequent decays (because D has hypercharge, the LEP2 bound of ~ 100 GeV on its mass applies [11]). The other important

production process is the pair productions of the exotic quark in association with the new $U(1)'$ gauge boson, $D\bar{D}Z'$. It turns out that this is the only way the new gauge boson Z' can be produced on-shell at LHC because of its very suppressed or vanishing couplings to the SM particles in this model. In the following subsections we discuss the signals from the $D\bar{D}$ production. We also discuss the couplings of the extra gauge boson Z' with the SM particles.

A. Signals from $D\bar{D}$ Productions

The heavy exotic quarks being colored particles will be produced copiously at the LHC through strong interactions. The major contribution, as in the case of top quarks, would come from the gluon-induced subprocess ($\sim 80\%$). In Fig. 1 we plot the pair production cross section for the process

$$pp \rightarrow D\bar{D} \quad (16)$$

at LHC for two different center-of-mass energies (7 TeV and 14 TeV).

The figure clearly shows that one can have quite large production cross section for such an exotic quark at the LHC and its signals should be observable through its decay products. We also include the most updated Tevatron bounds obtained by the Collider Detector at Fermilab (CDF) using 4.8 fb^{-1} data which rules out such exotic

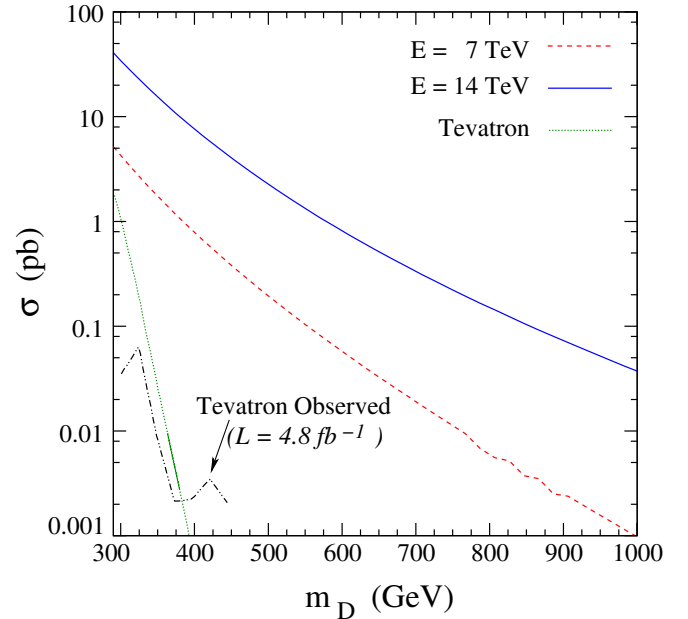


FIG. 1 (color online). Pair production cross section for the exotic quarks at LHC as a function of its mass (m_D). We use the CTEQ6L1 parton distribution functions (PDF) [22] for the protons. We have set the scale $Q^2 = m_D^2$. We also include the cross section at the Tevatron experiment along with the observed data by CDF [12], excluding such an exotic quark of mass ≤ 385 GeV at 95% C.L.

TABLE I. The effective coupling of the exotic D quark with the other particles in the model. The electromagnetic coupling with photon and the strong coupling with gluon is the same as any down-type quarks in the SM. Couplings are of the form $K\gamma^\mu(c_V - c_A\gamma^5)$ and $K(c_S - c_P\gamma^5)$. Note that we have put $V_{tb} = 1$.

	K	c_V	c_A
$\bar{b}bZ_\mu$	$\frac{e}{12\sin 2\theta_W}$	$4\cos 2\theta_W + 3\cos 2\theta_L - 1$	$3(1 + \cos 2\theta_L)$
$\bar{D}DZ_\mu$	$\frac{e}{12\sin 2\theta_W}$	$4\cos 2\theta_W - 3\cos 2\theta_L - 1$	$3(1 - \cos 2\theta_L)$
$\bar{b}DZ_\mu$	$\frac{-e}{4}\frac{\sin 2\theta_L}{\sin 2\theta_W}$	1	1
$\bar{b}bZ'_\mu$	$\frac{g_{Z'}}{4}Y_{Z'}$	$\cos 2\theta_L + \cos 2\theta_R - 2$	$\cos 2\theta_L - \cos 2\theta_R$
$\bar{D}DZ'_\mu$	$\frac{-g_{Z'}}{4}Y_{Z'}$	$\cos 2\theta_L + \cos 2\theta_R + 2$	$\cos 2\theta_L - \cos 2\theta_R$
$\bar{b}DZ'_\mu$	$\frac{-g_{Z'}}{4}Y_{Z'}$	$\sin 2\theta_L + \sin 2\theta_R$	$\sin 2\theta_L - \sin 2\theta_R$
$\bar{t}bW_\mu^+$	$\frac{-e\cos\theta_L}{2\sqrt{2}\sin\theta_W}$	1	1
$\bar{t}DW_\mu^+$	$\frac{e\sin\theta_L}{2\sqrt{2}\sin\theta_W}$	1	1
	K	c_S	c_P
$\bar{b}b\phi_H$	$\frac{\cos\theta_R}{\sqrt{2}}$	$y_b\cos\beta\cos\theta_L + y_{Db}\sin\beta\sin\theta_L$	0
$\bar{D}D\phi_H$	$\frac{\sin\theta_R}{\sqrt{2}}$	$y_b\cos\beta\sin\theta_L - y_{Db}\sin\beta\cos\theta_L$	0
$\bar{b}b\phi_S$	$\frac{-\cos\theta_R}{\sqrt{2}}$	$y_b\sin\beta\cos\theta_L - y_{Db}\cos\beta\sin\theta_L$	0
$\bar{D}D\phi_S$	$\frac{-\sin\theta_R}{\sqrt{2}}$	$y_b\sin\beta\sin\theta_L + y_{Db}\cos\beta\cos\theta_L$	0
$\bar{b}D\phi_H$	$\frac{1}{2\sqrt{2}}$	$y_{Db}\sin\beta\cos(\theta_L + \theta_R) - y_b\cos\beta\sin(\theta_L + \theta_R)$	$y_{Db}\sin\beta\cos(\theta_L - \theta_R) - y_b\cos\beta\sin(\theta_L - \theta_R)$
$\bar{b}D\phi_S$	$\frac{1}{2\sqrt{2}}$	$y_{Db}\cos\beta\cos(\theta_L + \theta_R) + y_b\sin\beta\sin(\theta_L + \theta_R)$	$y_{Db}\cos\beta\cos(\theta_L - \theta_R) + y_b\sin\beta\sin(\theta_L - \theta_R)$

quarks with mass $m_D < 385$ GeV at 95% C.L. [12]. However, this bound assumes that the exotic bottom-type quark decays to tW with 100% probability. The bound gets diluted as the branching probability goes down, and in our case turns out to be ~ 300 GeV for our choice of model parameters. We have implemented the model into CALCHEP [13] to calculate the production cross sections as well as the two-body decays of the new particles in the model.

The heavy quark in the gauge eigenbasis couples directly to the Z' gauge boson through the $U(1)'$ charge, with the gauge coupling strength $g_{Z'}$. However, its decay is more dependent on the mixing parameters resulting from its mixing with the b quark, leading to a much richer phenomenology. The heavy exotic itself is a mixed mass eigenstate and we list its couplings to the other particles of the model in Table I.

The two-body decay width for D of mass m_D , in its rest frame can be written down as

$$\Gamma(D \rightarrow X_2 X_3) = \frac{1}{16\pi m_D} \lambda^{1/2}\left(1, \frac{m_{X_2}^2}{m_D^2}, \frac{m_{X_3}^2}{m_D^2}\right) |\overline{\mathcal{M}}|^2, \quad (17)$$

where the function $\lambda(x, y, z) = x^2 + y^2 + z^2 - 2(xy + yz + zx)$. Using the effective couplings given in Table I, one can write down the explicit decay amplitudes for the exotic quarks decaying into vector (V) and scalar (Φ) modes

$$\begin{aligned}
|\overline{\mathcal{M}}|^2(D \rightarrow fV) &= K^2 \left[3 \left(m_D^2 + m_f^2 - 2m_V^2 + \frac{(m_D^2 - m_f^2)^2}{m_V^2} \right) \right. \\
&\quad \left. \times (c_V^2 + c_A^2) - 18m_D m_f (c_V^2 - c_A^2) \right] \\
|\overline{\mathcal{M}}|^2(D \rightarrow f\Phi) &= K^2 [3(m_D^2 + m_f^2 - m_\Phi^2)(c_S^2 + c_P^2) \\
&\quad + 6m_D m_f (c_S^2 - c_P^2)]. \quad (18)
\end{aligned}$$

We can now estimate the decay probabilities of the heavy exotic D quark. To highlight distinct scenarios, we choose two different sets of input values for the free parameters as representative points in the model listed in Table II. Note that the input parameters for the model also affect some EW observables, e.g., the Z boson decay width or the mass

TABLE II. Representative points in the model parameter space and the relevant mass spectrum used in the analysis.

Parameters	I	II
$(\lambda_H, \lambda_S, \lambda_{HS})$	(0.11, 0.16, 0.005)	(0.2, 0.05, 0.1)
v_S	1000 GeV	800 GeV
Y_{Db}	0.15	0.05
m_{ϕ_H}	115 GeV	127 GeV
m_{ϕ_S}	566 GeV	268 GeV
$m_{Z'}$	1000 GeV	800 GeV

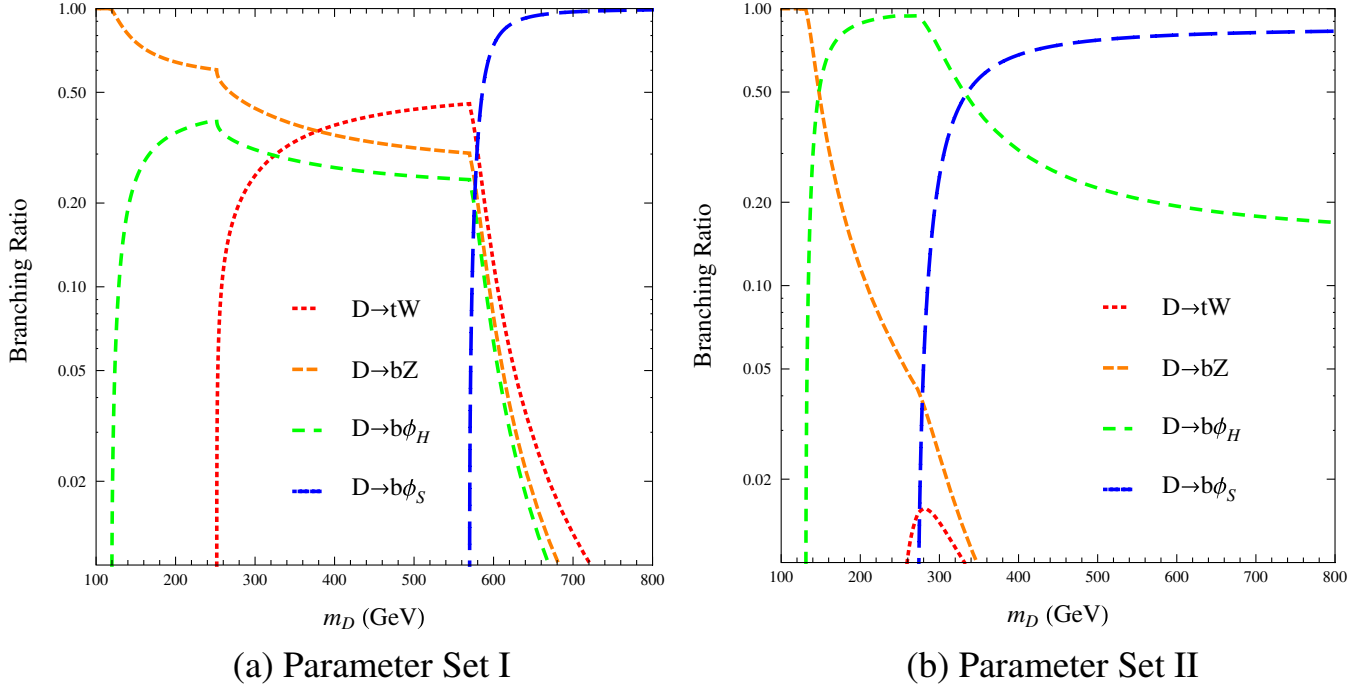


FIG. 2 (color online). Illustrating the decay probabilities of the D quark as a function of its mass m_D : (a) parameter set I and (b) parameter set II.

limits for Higgs boson and other heavy exotics that appear in our model. We have checked that the input parameters given in Table II are allowed and do not contradict any existing experimental bounds [11]. In Fig. 2 we present the decay branching ratios (BR) of the heavy quark D as a function of its mass (m_D) for the representative points I and II given in Table II. The curves in Fig. 2(a) represent point I

from Table II. When D is lighter than $m_t + M_W$ then it always decays to Zb through mixing if its coupling to the lighter Higgs boson is very suppressed. This would happen when the lighter scalar state is dominantly an $SU(2)$ doublet. The tW^- mode starts picking up and becomes comparable to the Zb mode for heavier D . tW^- is a common decay mode in fourth-generation models and theories with

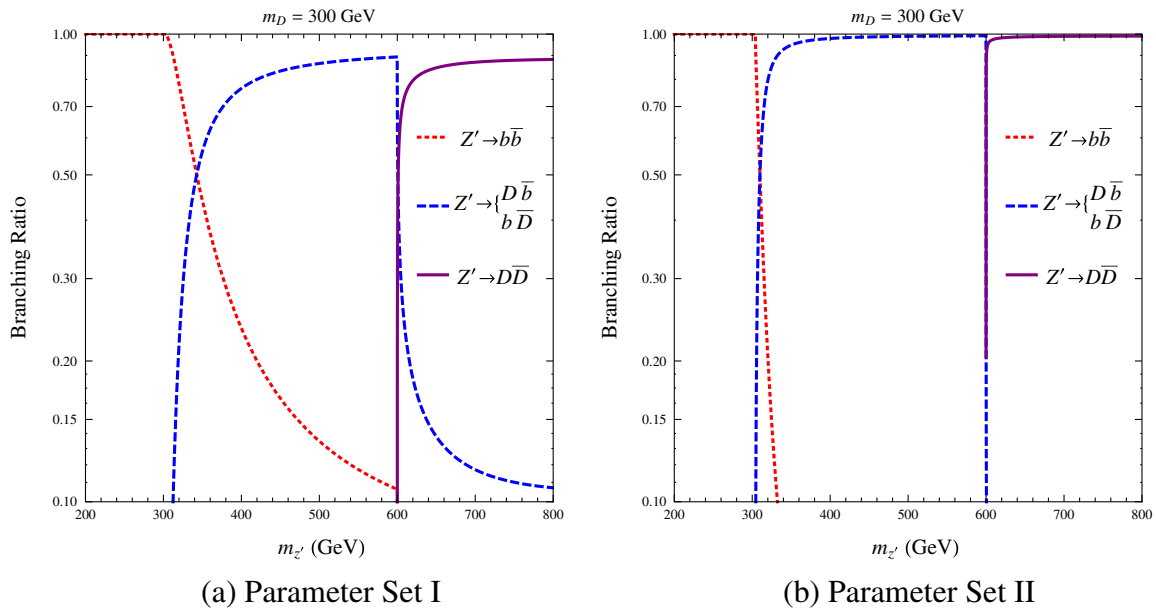


FIG. 3 (color online). Illustrating the decay probabilities of the Z' as a function of its mass $m_{Z'}$: (a) parameter set I and (b) parameter set II.

TABLE III. Branching ratios for various Higgs decay modes for parameter sets I and II.

Decays	Branching ratios		Decays	Branching ratios	
	I	II		I	II
$\phi_H \rightarrow b\bar{b}$	0.672	0.510	$\phi_S \rightarrow \phi_H \phi_H$	0.25	0.27
$\phi_H \rightarrow c\bar{c}$	0.031	0.024	$\phi_S \rightarrow W^+ W^-$	0.42	0.51
$\phi_H \rightarrow \tau^+ \tau^-$	0.093	0.072	$\phi_S \rightarrow ZZ$	0.20	0.22
$\phi_H \rightarrow gg$	0.104	0.096	$\phi_S \rightarrow t\bar{t}$	0.13	–
$\phi_H \rightarrow WW^*$	0.088	0.266			

top or bottom partners as studied in Ref. [14] and results in multilepton signals.

The curves in Fig. 2(b) represent point II, where the choice of parameters gives a very suppressed mixing angle θ_L . The couplings of Zb and tW to the exotic quark are proportional to $\sin\theta_L$ and hence also get suppressed. As soon as the scalar modes become kinematically accessible, they completely dominate the decay properties of the exotic quark.

We find if the Z' boson is light, then as soon as the $D \rightarrow Z'b$ mode opens up, the remaining modes drop out very quickly for point I while for point II it becomes comparable to the scalar mode for very large mass m_D . It is worth pointing out here that the dominant decay mode for the Z' when $m_{Z'} < m_D$ is to $b\bar{b}$ with 100% branching probability. However for Z' heavier than D the dominant decay of Z' is to $D\bar{b}$ and $\bar{D}b$ as shown in Fig. 3; and as soon as $m_{Z'} > 2m_D$ it decays dominantly to $\bar{D}D$ with maximum probability. We should also point out that the Z' phenomenology in our model is quite different from other models with $U(1)$ extension of the SM. As there exists no coupling between any SM fermion pair (other than b quark) or with the EW gauge bosons (no kinetic mixing), it is not possible to produce this particle directly through exchange of SM particles at LEP or Tevatron and so the strong constraints that exist on the mass of similar Z' exotics through the effective four-fermion operators [15] do not apply in our case and neither do the search limits from the Tevatron experiments [16]. Thus the Z' in our model can be light but remain invisible in the existing experimental data. We will however not discuss the Z' signals any further and only focus on the signals arising from the production of the

exotic D quarks in the model. To understand the full decay chain of the D quark to final state particles we also list the decay probabilities of the scalars ϕ_H and ϕ_S in Table III for the two representative points I and II.

Thus the above decay patterns suggest that one can have the following interesting final states from the decay of the exotic quarks:

$$pp \rightarrow D\bar{D}$$

$$\rightarrow \begin{cases} (Zb)(Z\bar{b}) & \Rightarrow b\bar{b} + 2Z \\ (tW^-)(\bar{t}W^+) & \Rightarrow t\bar{t} + W^+ W^- \\ (tW^-)(Z\bar{b}) & \Rightarrow t\bar{b} + ZW^- \\ (\phi_H b)(\phi_H \bar{b}) & \Rightarrow b\bar{b} + 2\phi_H \rightarrow b\bar{b} + 2(W^+ W^-) \\ (\phi_H b)(\phi_H \bar{b}) & \Rightarrow b\bar{b} + 2\phi_H \rightarrow 3(b\bar{b}) \\ (\phi_H b)(\phi_S \bar{b}) & \Rightarrow b\bar{b} + \phi_H \phi_S \rightarrow b\bar{b} + 3\phi_H \rightarrow 4(b\bar{b}) \\ (\phi_S b)(\phi_S \bar{b}) & \Rightarrow b\bar{b} + 2\phi_S \rightarrow b\bar{b} + 4\phi_H \rightarrow 5(b\bar{b}) \end{cases}, \quad (19)$$

where the first four suggest multilepton and multijet final states with two or more b jets, while the remaining give more exotic signatures like N b -jet final states where N can be as large as 10. Note that the above decays only illustrate some of the possible decay chains and we have not listed other possible combinations of the D decays which can also lead to similar final states.

In Table IV we list the probabilities for the decay modes for a few specific values of the D quark mass. We also show the corresponding cross sections for the pair production of these exotics at LHC for the center-of-mass energies of $\sqrt{s} = 7$ TeV and $\sqrt{s} = 14$ TeV. The decays suggest a large multiplicity of b quarks in the final state. It turns out that six- b final states for the signal is very promising.

TABLE IV. Cross sections and branching probabilities for specific mass values of D quark for the representative points I and II.

m_D	$\sigma(D\bar{D})$ (pb)		Branching ratios							
			$D \rightarrow tW$		$D \rightarrow bZ$		$D \rightarrow b\phi_H$		$D \rightarrow b\phi_S$	
	7 TeV	14 TeV	I	II	I	II	I	II	I	II
250	12.15	87.760	0	0	0.603	0.055	0.397	0.945	0	0
300	4.265	34.368	0.251	0.014	0.438	0.024	0.311	0.715	0	0.247
400	0.791	7.692	0.381	0.005	0.351	0.005	0.268	0.308	0	0.681
500	0.194	2.270	0.434	0.003	0.316	0.002	0.250	0.225	0	0.770
600	0.059	0.820	0.121	0.002	0.078	0.001	0.063	0.194	0.738	0.803

However, there exists no estimate for this final state in the literature, arising from the SM. We present below a leading-order (LO) estimate of the cross section for the six- b SM background from QCD for LHC energies.

B. Calculation of six- b final states from QCD

The six- b final state is interesting, independent of our particular model. The presence of six b jets allow the jets to be tagged. All other six-jet final states involve mixtures of light quarks and gluons, and one cannot separate light jets from gluon jets. Therefore the six- b final state itself presents an interesting test of QCD. Furthermore by computing the full matrix element, we can test the validity of the differential cross section by looking at differential observables. While we only use the six- b cross section as a background to our signal, this is the first time such a six- b cross section in QCD has been estimated, and this also is an important result of this paper.

Any six-particle final state process is a challenge to compute. The phase space is 14 dimensional (excluding the two additional integrations over the initial parton fluxes). There are thousands of diagrams, and thousands of distinct color configurations. While six-jet final states are produced every day by Monte Carlo generators such as PYTHIA [17] and HERWIG [18], the mechanism they use creates additional jets from an initial $2 \rightarrow 2$ or $2 \rightarrow 3$ process via a showering procedure that resums leading logarithms, splitting an extra gluon from the hard final state partons.

As is well known, the showering procedure cannot produce the correct correlations among three or more hard jets, nor can it compute the total cross section for three or more hard jets. It assumes that each parton is independent of all the others. For a single radiation it is strictly correct in the limit that the extra radiation is soft and/or collinear with the initial parton. However for multiple radiations it ignores the QCD connection among the radiations, assuming that each radiation factorizes from the others. There is also quantum mechanical interference in different radiations which result in the same final state that is ignored.

This means that one should not examine in detail observables such as the angles between jets, invariant masses of jet pairs, or the thrust, when one hard jet came from the showering procedure. This showering technique is however extremely useful as long as one is not sensitive to the

details of correlations in the differential cross section, as this method is computationally simpler than a full matrix element calculation.

Therefore to have an accurate Monte Carlo with six jets in the final state, one must compute the full matrix element, which automatically includes all color flows and interference. This is accurate to approximately the 10% level, at which point next-to-leading order loop corrections become important. Note that due to the b quark mass, there is no soft radiation which benefits significantly from the usual Sudakov logarithm resummation. The b mass acts as a regulator, relegating this cross section strictly into the “hard” regime, in which the matrix element is valid. Even if all six b ’s are at rest, the energy in the final state is $6m_b$, and any virtual gluon must have a virtuality $q^2 \sim m_b^2$.

For this calculation we have chosen the tool MADEVENT, which is an event generator built upon the matrix element generator MADGRAPH [19]. We have modified these tools to be able to cope with thousands of diagrams and thousands of color configurations. Computing five-particle and six-particle final state QCD processes has a number of challenges, all of which are technical rather than physics-based. MADEVENT is in principle capable of computing any process with any number of final state particles, however several internal restrictions caused previous versions to fail on processes such as six b ’s. Although the program is equipped to handle the large color configurations, some input/output statements restrict this computation to a smaller number of color flow configurations which makes it incapable of calculating the six- b final states at a hadron machine like the LHC. These have been repaired to calculate the six- b cross section in the SM from QCD at the LHC.

C. Signal and background analysis

A simple minded estimate of the cross section using $\sigma \times \text{BR}$ shows that the final states which would be of interest at the LHC would involve at least two b jets in the final state. Besides the two hard b jets, one expects charged leptons in the final state coming from the decays of the weak gauge bosons. It is also worth noting that when the D quark decays to the Higgs bosons, one would get a large multiplicity of b jets in the final states as the Higgs with $M_{\phi_H} < 2M_W$ dominantly decays to b jets.

TABLE V. Illustrating the final state cross sections after the decay of D quarks. All cross sections are in units of femtobarn (fb).

Final States	$m_D = 300 \text{ GeV}$				$m_D = 500 \text{ GeV}$			
	$\sqrt{s} = 7 \text{ TeV}$		$\sqrt{s} = 14 \text{ TeV}$		$\sqrt{s} = 7 \text{ TeV}$		$\sqrt{s} = 14 \text{ TeV}$	
	I	II	I	II	I	II	I	II
$6b + X$	181.92	718.79	1394.32	5521.23	4.94	10.79	531.10	115.02
$nb + \ell + X \ (n \geq 2)$	452.50	188.22	3559.94	1465.42	32.17	27.88	43.37	313.96
$2b + 2\ell + X$	146.53	14.95	1127.15	117.08	8.71	6.61	11.77	75.56
$4b + 2\ell + X$	51.07	24.36	384.24	183.73	1.85	2.06	20.33	22.20

To select our events for the final states given in Table V, we have imposed the following kinematic cuts:

- (i) All the b jets must have a $p_T^b > 20$ GeV and lie within the rapidity gap of $|\eta^b| < 3.0$.
- (ii) All charged leptons ($\ell = e, \mu$) must have a $p_T^\ell > 20$ GeV and lie within the rapidity gap of $|\eta^\ell| < 2.5$.
- (iii) The final states also must satisfy $\Delta R_{bb} > 0.7$, $\Delta R_{b\ell} > 0.4$, and $\Delta R_{\ell\ell} > 0.2$ where $\Delta R_{ij} = \sqrt{(\Delta\eta_{ij})^2 + (\Delta\phi_{ij})^2}$.
- (iv) All b -jet pairs must have a minimum invariant mass $M_{bb} > 10$ GeV.

In Table V we present the cross sections for the signal for two different mass values of the exotic D quark after passing through the above mentioned kinematic cuts. As expected, the favored final states are dependent on the high b -jet multiplicity. At the hadron collider such as LHC, one favors final states with leptons. However b jets can also be triggered upon and identified and thus can prove to be useful in isolating new physics signals such as ours which involve at least two or more b jets. Looking at Table V we find that we get a good signal rate for the inclusive $6b + X$ final state. The SM background for multi- b final state is quite large [20]. However no estimate of a six- b final state exists in the literature, which we find relevant for our signal. We have used the MADGRAPH+MADEVENT package to estimate the leading-order partonic cross section for the six- b final state at LHC. With the above mentioned kinematic cuts, we find that for LHC energy of $\sqrt{s} = 14$ TeV, the SM background for a six- b final state is ~ 70 fb and falls to less than 10 fb for the $\sqrt{s} = 7$ TeV option. This implies that the signal in our model is much larger than the SM background even for larger mass values of the exotic D quark. The other signals which are worth looking for in this model is one or two charged leptons with varying b -jet multiplicities. We have listed the interesting ones in Table V. The final state with $2\ell + 4b + X$ also stands out against the SM background, where one gets the SM cross sections to be quite small as it is already α_{EW}/α_s suppressed compared to the six- b cross section. The SM background is much larger for the final states $\ell + nb + X$ where $n = 2$ and $2\ell + 2b + X$, where the significant SM background results from the $t\bar{t}$ production. For the final states $\ell + nb + X$ one can get rid of the huge $t\bar{t}$ background by demanding $n \geq 3$. This helps in improving the significance of the signal, even though we also lose a large fraction of the signal events in the process. For the other final state, we find that at leading order, at LHC with pp collision energy of $\sqrt{s} = 7$ TeV, the $2\ell + 2b + X$ SM background is ~ 3.3 pb. As the leading two b jets in our signal come from the decays of the heavy exotic D quark, we put a stronger p_T cut of 100 GeV. This reduces the signal by two-thirds. However the SM background is reduced by more than an order of magnitude, and becomes

~ 232 fb for $\sqrt{s} = 7$ TeV collisions while it is ~ 1.63 pb for $\sqrt{s} = 14$ TeV which does look promising for the signal with large enough luminosities at the LHC. We must point out that we have not incorporated any efficiency factors for our final state particles. Most notably, all numerical estimates involving b jets for signal as well as the SM background will have to be scaled with the b -tagging efficiency of around 50–60% expected at the LHC [21].

IV. SUMMARY AND CONCLUSIONS

In this work, we have proposed a new extension of the SM, by introducing a hidden $U(1)'$ symmetry. The difference with the previously studied $U(1)'$'s is that all the SM particles are singlets under our proposed new $U(1)'$, and hence hidden. Such a symmetry may be present at the TeV scale, and may manifest at the LHC giving new signals observable at the LHC. The model incorporates a new EW singlet Higgs, as well as new vectorlike charge $-1/3$ quarks. The new massive vector boson, Z' resulting from the broken hidden $U(1)'$ symmetry does not couple to the SM fields directly and may be hidden in all experimental data due to the very small mixing parameters making its production cross section too small to observe. However, if once produced, it can be observed through its decays. We leave its detailed phenomenological study for future work and have focused on the exotic colored fermion in the model. We have studied the pair productions of these new quarks and their subsequent decays. The dominant final states include multiple b jets with high p_T , or b jets plus charged leptons with high p_T and missing energy, and stand out beyond the SM background. The most distinctive final state signal is the six- b quark with high p_T and no missing energy. A lot of effort have been put in both for the ATLAS and CMS detectors to improve the b -tagging efficiency. So the calculation for this six- b final state is also of great importance in the SM, and has not yet been calculated. We have calculated this six- b signal in our model, and have also estimated the SM expectation using MADGRAPH and MADEVENT. We found that the signal in our model stands well above that expected from the SM.

ACKNOWLEDGMENTS

We thank T. Stelzer and F. Maltoni for several useful communications regarding the MADGRAPH and MADEVENT packages. S. Nandi thanks the Fermilab Theory Group for their warm hospitality and support from their Summer Visitor Program during the completion of this work. S. Nandi also thanks the CERN Theory Group for warm hospitality and support while part of this work was done. S.K. Rai would like to acknowledge and thank the Helsinki Institute of Physics for their support and hospitality. The work of B. N. G., S. N., and S. K. R. is supported in part by the United States Department of Energy, Grant Nos. DE-FG02-04ER41306 and DE-FG02-04ER46140.

- [1] H. Davoudiasl, R. Kitano, T. Li, and H. Murayama, *Phys. Lett. B* **609**, 117 (2005).
- [2] I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, and G.R. Dvali, *Phys. Lett. B* **436**, 257 (1998).
- [3] I. Gogoladze, Y. Mimura, and S. Nandi, *Phys. Lett. B* **562**, 307 (2003).
- [4] J.C. Pati and A. Salam, *Phys. Rev. D* **10**, 275 (1974); **11**, 703(E) (1975).
- [5] See, for example, P. Langacker, [arXiv:0911.4294](#); J.L. Hewett and T.G. Rizzo, *Phys. Rep.* **183**, 193 (1989).
- [6] D.J. Muller and S. Nandi, *Phys. Lett. B* **383**, 345 (1996); E. Malkawi, T.M.P. Tait, and C.P. Yuan, *Phys. Lett. B* **385**, 304 (1996); R.S. Chivukula, E.H. Simmons, and J. Terning, *Phys. Rev. D* **53**, 5258 (1996).
- [7] T. Han, P. Langacker, and B. McElrath, *Phys. Rev. D* **70**, 115006 (2004).
- [8] G. Bertone, D. Hooper, and J. Silk, *Phys. Rep.* **405**, 279 (2005).
- [9] For previous work on hidden non-Abelian gauge symmetries, see T.j. Li and S. Nandi, *Phys. Lett. B* **617**, 112 (2005); S. Gopalakrishna, S. Jung, and J.D. Wells, *Phys. Rev. D* **78**, 055002 (2008).
- [10] V. Barger, H.E. Logan, and G. Shaughnessy, *Phys. Rev. D* **79**, 115018 (2009); G. Bhattacharyya, G.C. Branco, and S. Nandi, *Phys. Rev. D* **77**, 117701 (2008).
- [11] C. Amsler *et al.* (Particle Data Group), *Phys. Lett. B* **667**, 1 (2008).
- [12] L. Scodellaro (CDF Collaboration and D0 Collaboration), in *Proceedings of the International Conference on High Energy Physics, Paris, 2010* (unpublished).
- [13] A. Pukhov, [arXiv:hep-ph/0412191](#).
- [14] R. Contino and G. Servant, *J. High Energy Phys.* **06** (2008) 026; J.A. Aguilar-Saavedra, *J. High Energy Phys.* **11** (2009) 030; T. Han, K. Zurek, M. Strassler, and Z. Si, *J. High Energy Phys.* **07** (2008) 008; W.F. Chang, J.N. Ng, and J.M.S. Wu, *Phys. Rev. D* **75**, 115016 (2007).
- [15] J. Erler, P. Langacker, S. Munir, and E.R. Pena, *J. High Energy Phys.* **08** (2009) 017.
- [16] T. Aaltonen *et al.* (CDF Collaboration), *Phys. Rev. Lett.* **102**, 091805 (2009).
- [17] T. Sjostrand, S. Mrenna, and P.Z. Skands, *J. High Energy Phys.* **05** (2006) 026.
- [18] G. Corcella *et al.*, [arXiv:hep-ph/0210213](#).
- [19] T. Stelzer and W.F. Long, *Comput. Phys. Commun.* **81**, 357 (1994); F. Maltoni and T. Stelzer, *J. High Energy Phys.* **02** (2003) 027.
- [20] H. Baer, V. Barger, A. Lessa, and X. Tata, *J. High Energy Phys.* **09** (2009) 063.
- [21] I.R. Tomalin (CMS Collaboration), *J. Phys. Conf. Ser.* **110**, 092033 (2008).
- [22] J. Pumplin, A. Belyaev, J. Huston, D. Stump, and W.K. Tung, *J. High Energy Phys.* **02** (2006) 032.