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Realistic models with a light U(1) gauge boson coupled to baryon number

Christopher D. Carone and Hitoshi Murayama*

Theoretical Physics Group, Lawrence Berkeley Laboratory, University of California, Berkeley, California 94720

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We recently showed that a new gauge boson γ_B coupling only to baryon number is phenomenologically allowed, even if $m_B < m_Z$ and $\alpha_B \approx 0.2$. In our previous work we assumed that kinetic mixing between the baryon number and hypercharge gauge bosons (via an ${F_B}^{\mu\nu}{F_{\mu\nu}}^Y$ term) was small enough to evade constraints from precision electroweak measurements. In this paper we propose a class of models in which this term is naturally absent above the electroweak scale. We show that the generation of a mixing term through radiative corrections in the low-energy effective theory does not lead to a conflict with precision electroweak measurements and may provide a leptonic signal for models of this type at an upgraded Fermilab Tevatron.

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I. INTRODUCTION

In a recent paper, we considered the phenomenology of a light U(1) gauge boson γ_B that couples only to baryon number [1]. We assumed that the new U(1) gauge symmetry is spontaneously broken, and that the γ_B mass m_B is smaller than m_Z. Nevertheless, we showed that this new gauge boson could remain undetected, even if the coupling α_B were comparable to $\alpha_{\rm strong}$ [1, 2]. Since the γ_B boson couples only to quarks, any process that is relevant to γ_B detection also has a significant contribution from QCD. Thus, a typical γ_B boson with $m_B = 50$ GeV and $\alpha_B = 0.1$ can remain undetected by "hiding" in the large QCD background. Since the γ_B boson couples only to quarks, it is difficult to detect, just like the more familiar example of a light gluino in supersymmetric models [3].

One of the assumptions in our original analysis [1] was that mixing between the γ_B and electroweak gauge bosons was negligible. Mass mixing is not present because we assume that there are no Higgs bosons that carry both baryon number and electroweak quantum numbers. However, there is a possible off-diagonal kinetic term that mixes the $\mathrm{U}(1)_B$ and $\mathrm{U}(1)_Y$ gauge fields:

$$\mathcal{L}_{\rm kin} = -\frac{1}{4}\left(F_Y^{\mu\nu}F_{\mu\nu}^Y + 2c\,F_B^{\mu\nu}F_{\mu\nu}^Y + F_B^{\mu\nu}F_{\mu\nu}^B\right). \quad (1.1)$$

Here the $F^{\mu\nu}$ are gauge field strength tensors, and c is an undetermined coupling constant. Clearly, c must be quite small so that the kinetic mixing does not conflict with precision electroweak measurements. Although the phenomenology of the γ_B is specified within the three-dimensional parameter space m_B-α_B-c, any realistic model must reside within the narrow region $|c| < c_0$, where $c_0 \ll 1$ can be determined from the precision electroweak constraints. Thus, in our previous work [1], we described the γ_B phenomenology in terms of an effectively two-dimensional parameter space, the m_B-α_B plane at $c \approx 0$.

The natural question that remains to be answered is whether there are any models in which c is naturally small enough to satisfy the experimental constraints. Our previous results would be greatly undermined if they were relevant only to models in which the coupling c required fine-tuning at the electroweak scale. In this paper, we will describe a class of models in which this kinetic mixing term is absent above some scale Λ that we assume is not much greater than the top quark mass. Below Λ, a kinetic mixing term is generated only through radiative corrections, so that $c(\Lambda) = 0$ but $c(\mu) \neq 0$ for $\mu < \Lambda$. We will show that if 200 GeV $\lesssim \Lambda \lesssim$ 1.3 TeV, $c(\mu)$ never becomes large enough in the low-energy theory to conflict with precision electroweak measurements, even when α_B is as large as 0.1. We then present a model that satisfies this boundary condition. In addition, we show that a mixing term small enough to satisfy the current experimental constraints can nonetheless provide us with a possible signal for the γ_B in the Drell-Yan dilepton differential cross sections at hadron colliders. This signal could be within the reach of the Fermilab Tevatron with the main injector and a luminosity upgrade.

The paper is organized as follows. In the next section we discuss the phenomenological constraints on the kinetic mixing term from precision electroweak measurements. We show that these constraints can be satisfied if the scale Λ at which the mixing vanishes is just above the electroweak scale. In Sec. III, we present a model with gauged baryon number in which the kinetic mixing term is naturally absent above Λ. In Sec. IV we discuss a leptonic signature of the γ_B in Drell-Yan dilepton production at hadron colliders. In the final section we summarize our conclusions.

II. MIXING CONSTRAINTS

To study the effects of the kinetic mixing term, we could redefine the gauge fields so that the kinetic terms in

*On leave of absence from Department of Physics, Tohoku University, Sendai 980, Japan.

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the new basis are diagonal and conventionally normalized [4]. However, since we know *a priori* that the coupling c must be much less than 1, it is more convenient to treat the mixing term in Eq. (1.1) as a new perturbative interaction. The Feynman rule corresponding to the γ_B-photon vertex is

$$-ic_\gamma \cos\theta_w (p^2 g^{\mu\nu} - p^\mu p^\nu) \ , \tag{2.1}$$

and for the γ_B-Z^0 vertex is

$$ic_Z \sin\theta_w (p^2 g^{\mu\nu} - p^\mu p^\nu). \tag{2.2}$$

Note that $c_\gamma = c_Z = c$ above the electroweak scale, but c_γ and c_Z run differently in the low-energy effective theory below $m_{\rm top}$. This can be seen in Fig. 1. If we assume that both c_Z and c_γ vanish at the scale $\Lambda \approx m_{\rm top}$, we can run the couplings down to lower energies to determine the magnitude of mixing that is characteristic of purely radiative effects. This will give us a useful point of reference when we consider the relevant experimental constraints. If Λ were less than $m_{\rm top}$, the new fermions that are associated with this scale (cf. Sec. III) would have already been seen in accelerator experiments. The extent to which Λ can be significantly larger than $m_{\rm top}$ will be considered later in this section.

The couplings c_Z and c_γ are renormalized by the one quark-loop diagrams that connect the γ_B to either the photon or Z. From these diagrams we obtain the one-loop renormalization group equations

$$\mu\frac{\partial}{\partial\mu}c_\gamma(\mu) = -\frac{2}{9\pi}\frac{\sqrt{\alpha_B\alpha}}{c_w}\,[2N_u - N_d] \tag{2.3}$$

and

$$\mu\frac{\partial}{\partial\mu}c_Z(\mu) = -\frac{1}{18\pi}\frac{\sqrt{\alpha_B\alpha}}{s_w^2 c_w}[3(N_d - N_u) + 4(2N_u - N_d)s_w^2], \tag{2.4}$$

where s_w (c_w) $= \sin\theta_w$ ($\cos\theta_w$), and α is the fine structure constant. N_u and N_d are the number of light quark flavors with charge 2/3 and $-1/3$, respectively. For $\Lambda = m_{\rm top}$ and $\alpha_B = 0.1$, we find

$$c_\gamma(m_Z) \approx 1.5\times10^{-3}, \quad c_Z(m_Z) \approx 6.2\times10^{-3},$$
$$c_\gamma(m_b) \approx 8.4\times10^{-3}, \quad c_Z(m_b) \approx 3.5\times10^{-2},$$

where m_b is the bottom quark mass. As we cross the various quark mass thresholds below m_b, the rate at which c_γ and c_Z run becomes progressively smaller, until the running effectively stops below the hadronic scale ≈ 1 GeV. It is our first job to determine whether the estimates above are consistent with experimental limits on c_γ and c_Z. We will compare our results to the two standard deviation experimental uncertainty in each observable that is affected by the mixing. While strictly speaking it would be more accurate to do a global fit to the data while allowing c_Z, c_γ, α_B and the couplings of the standard model to vary, our approach will be sufficient to determine roughly whether or not the scenario we propose is allowed. Afterwards, we will return to the issue of the running and determine an upper bound on the scale Λ.

A. γ_B-Z mixing

The most significant constraints on $c_Z(m_Z)$ are shown in Fig. 2. We have considered the effects of the γ_B-Z mixing on the following experimental observables: the Z mass, hadronic width, and forward-backward asymmetries, and the neutral current νN and eN deep inelastic scattering cross sections. We now consider each of these in turn.

Z mass. We determine the shift in the Z mass by computing the shift in the real part of the pole in neutral gauge boson propagator. Thus, we set

$$\det\Gamma^{(2)}(p^2) = 0, \tag{2.5}$$

where

$$\Gamma^{(2)}(p^2) = \begin{pmatrix} p^2 - m_Z^2 + im_Z\Gamma_Z & -c_Z s_w p^2 \\ -c_Z s_w p^2 & p^2 - m_B^2 + im_B\Gamma_B \end{pmatrix}. \tag{2.6}$$

We find

$$\frac{\Delta m_Z}{m_Z} \approx 0.116\, c_Z^2 \frac{m_Z^2}{m_Z^2 - m_B^2}. \tag{2.7}$$

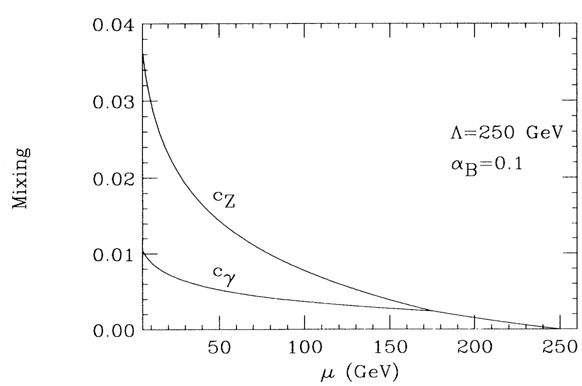

FIG. 1. Running of c_γ and c_Z, assuming $c_\gamma(\Lambda) = c_Z(\Lambda) = 0$ at $\Lambda = 250$ GeV.

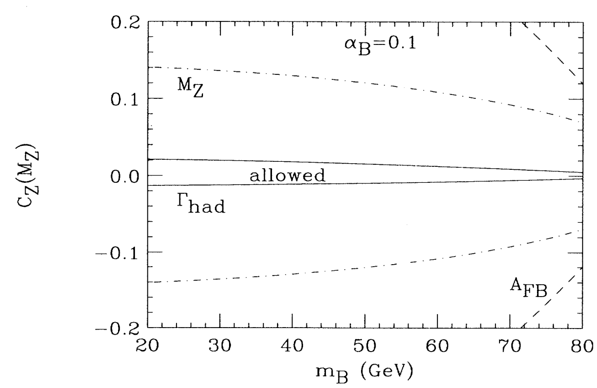

FIG. 2. Constraints on $c_Z(m_Z)$ from the two standard deviations of the experimental uncertainties in the Z mass, hadronic width, and $Z \to b\bar{b}$ forward-backward asymmetry.

Note that the effect of γ_B-γ mixing appears at $O(c_Z^2 c_\gamma^2)$, and is negligible. This expression is valid provided that m_B is not too close to m_Z. We have checked that this approximation is accurate if $|m_B - m_Z| \gtrsim 10$ GeV, which holds over range of m_B of interest to us in this paper. Since m_Z is taken as an input to determine other electroweak parameters, we require that the shift in m_Z does not spoil the consistency between the value of $\sin^2\hat{\theta}_w$ determined from the Z decay asymmetries (which we call s^2 below), and the value extracted from deep inelastic scattering data. The shift in m_Z corresponding to the uncertainty Δs^2 is given by

$$\frac{\Delta m_Z}{m_Z} = -\frac{1-2s^2}{2s^2c^2}\Delta s^2, \tag{2.8}$$

where $s^2 = 0.2317 \pm 0.0008$ [5]. Thus, we find $\Delta m_Z/m_Z < 2.4\times10^{-3}$, requiring that the shift in $\sin^2\hat{\theta}$ is no more than a two standard deviation effect. The contour corresponding to this bound is plotted in Fig. 2.

Z hadronic width. In Ref. [1] we computed the contribution to the Z hadronic width from (i) direct γ_B production $Z \to q\bar{q}\gamma_B$, and (ii) the $Zq\bar{q}$ vertex correction. There is an additional contribution to the hadronic width from the Z-γ_B mixing that is given by

$$\frac{\Delta\Gamma_{\rm had}}{\Gamma_{\rm had}} \approx -1.194\, c_Z\sqrt{\alpha_B}\frac{m_Z^2}{m_Z^2-m_B^2}. \tag{2.9}$$

Given that the uncertainty in the Z hadronic width is 0.6% at two standard deviations [5], we obtain the contour shown in Fig. 2. Notice that the constraint that we obtain is weaker for positive c_Z due to cancellation between the contributions discussed in Ref. [1] and the new contribution given in Eq. (2.9). For the typical values of $c_Z(m_Z)$ presented at the beginning of Sec. II, this cancellation is almost exact. In this case, we do not expect a large enough effect to account for the anomalously high value of $\alpha_{\rm QCD}$ reported by the CERN e^+e^- collider LEP, in contrast with the claim made in Ref. [2]. In any case, the hadronic width places the tightest constraint on $c_Z(m_Z)$, roughly $|c_Z(m_Z)| \lesssim 0.02$.

Forward-backward asymmetries. The γ_B-Z mixing term has the effect of slightly shifting the vector coupling of the Z to quarks. Thus, there is a new contribution to the forward-backward asymmetry $A_{\rm FB}^{(0,q)}$ in Z decay to $q\bar{q}$. Since the experimental uncertainty is smallest for $q=b$, we use the two standard deviation uncertainty in $A_{\rm FB}^{(0,b)}$ to constrain our model. We find that the new contribution is given by

$$\Delta A_{\rm FB}^{(0,b)} \approx -0.159\, c_Z\sqrt{\alpha_B}\frac{m_Z^2}{m_Z^2-m_B^2}, \tag{2.10}$$

while the measured value is 0.107 ± 0.013 [5]. The resulting bound is shown in Fig. 2. Notice that this provides a weaker constraint than those we obtained from consideration of the Z mass and width.

Deep inelastic scattering. The constraints on c_Z from deep inelastic νN scattering and from parity-violating eN scattering are much weaker than the other constraints that we have discussed and are not shown in Fig. 2. Deep inelastic νN scattering can be described in terms of the parameters $\epsilon_{L(R)}$, defined by the effective four-fermion interaction [5]

$$\begin{aligned} -\mathcal{L}^{\nu N} = \frac{G_F}{\sqrt{2}}&\bar{\nu}\gamma^\mu(1-\gamma^5)\nu \\ &\times\sum_i[\epsilon_L(i)\overline{q_i}\gamma_\mu(1-\gamma^5)q_i \\ &+\epsilon_R(i)\overline{q_i}\gamma_\mu(1+\gamma^5)q_i], \end{aligned} \tag{2.11}$$

where the sum is over quark flavors. We find that the contribution to the ϵ parameters from the γ_B-Z mixing is given by

$$\Delta\epsilon_{L(R)} \approx 0.766\, c_Z\sqrt{\alpha_B}\frac{q^2}{q^2-m_B^2}, \tag{2.12}$$

where $-q^2 \approx 20$ GeV^2 is a typical squared momentum transfer. The most accurately measured ϵ parameter is $\epsilon_L(d) = -0.438 \pm 0.012$ [5]. To demonstrate that the uncertainty in $\epsilon_L(d)$ provides only a weak constraint, we evaluate Eq. (2.12) for $m_B \approx 50$ GeV and $\alpha_B \approx 0.1$. We obtain the bound $|c_Z(q^2)| < 12.5$, at two standard deviations, which gives us $|c_Z(m_Z)| < 12.5$, because the contribution from the running is small. This is a much weaker constraint than the others that we have considered.

Parity-violating eN scattering can be described in terms of two other parameters C_1 and C_2 defined by the effective four-fermion interaction [5]

$$\begin{aligned} -\mathcal{L}^{eN} = -\frac{G_F}{\sqrt{2}}\sum_i&[C_{1i}\bar{e}\gamma_\mu\gamma^5 e\overline{q_i}\gamma^\mu q_i \\ &+C_{2i}\bar{e}\gamma_\mu e\overline{q_i}\gamma_\mu\gamma^5 q_i]. \end{aligned} \tag{2.13}$$

The γ_B-Z mixing contributes only to the parameter C_{1i}:

$$\Delta C_{1i} = -1.533\, c_Z\sqrt{\alpha_B}\frac{q^2}{q^2-m_B^2}. \tag{2.14}$$

The parameter measured with the least experimental uncertainty is $C_{1d} = 0.359 \pm 0.041$ [5]. If we again assume that $m_B = 50$ GeV and $\alpha_B = 0.1$, then the bound on $c_Z(m_Z)$ following from the two standard deviation uncertainty in C_{1d} is $|c_Z(m_Z)| < 21.2$. This is even weaker than the constraint we obtained from νN scattering. Note that there are no further constraints on ΔC_{1i} from the measurements of atomic parity violation because this process involves zero momentum transfer, where the kinetic mixing vanishes.

What we have seen is that the Z-pole observables place the tightest constraints on the mixing parameter c_Z, while deep inelastic scattering measurements do not provide any further constraints. Thus, if $|c_Z(M_Z)| < 0.02$, we are not likely to encounter any problems with the precision electroweak measurements that we have considered in this section.

B. γ_B-γ mixing

The coupling c_γ has its most significant effect on a different set of observables. Below we consider the effect of the γ_B-γ mixing on the cross section for $e^+e^- \to$ hadrons, and on the anomalous magnetic moments of the electron and muon.

$e^+e^- \to$ hadrons. The most important constraint on c_γ comes from the additional contribution to R, the ratio $\sigma(e^+e^- \to \text{hadrons})/\sigma(e^+e^- \to \mu^+\mu^-)$. We find

$$\frac{\Delta R}{R} = -1.8\, c_\gamma \sqrt{\alpha_B}\, \frac{s(s-m_B^2)}{(s-m_B^2)^2 + m_B^2\Gamma_B^2} + 8.9\, c_\gamma^2 \alpha_B\, \frac{s^2}{(s-m_B^2)^2 + m_B^2\Gamma_B^2}, \tag{2.15}$$

where $\sqrt{s}$ is the center of mass energy, and $\Gamma_B = \frac{1}{9} N_F \alpha_B m_B$ is the γ_B width, with $N_F = 5$ for the range of m_B of interest. Notice that the nonstandard contribution to R is maximized only in the vicinity of $s \approx m_B^2$. For any m_B of interest, we can constrain c_γ by considering the two standard deviation uncertainty in the value of R measured at $\sqrt{s} \approx m_B$. The results are shown in Fig. 3, based on the cumulative data on R taken at various values of $\sqrt{s}$ and compiled by the Particle Data Group [5]. Since there are values of $\sqrt{s}$ that have not been studied, the constraints on c_γ are strongest when m_B happens to coincide with a value of $\sqrt{s}$ at which there is an experimental data point available. Roughly speaking, the allowed region of Fig. 3 corresponds to $|c_\gamma(m_B)| < 0.01$. However, it is clear that the constraint can be significantly weaker if m_B happens to lie at a point where less data are available. Note that we have obtained strong constraints from (2.15) due to the resonant behavior of the amplitude for s-channel γ_B exchange; as we have seen earlier we obtain weaker constraints from t-channel processes (e.g., at the DESY ep collider HERA).

Anomalous magnetic moments. At the very lowest energies, we can constrain c_γ by the effect of the mixing on the anomalous magnetic moments of the electron and muon. Since this provides a much weaker constraint than the one we obtained from R, we do not show the result in Fig. 3. We find that the nonstandard contribution to the anomalous magnetic moment $a = (g-2)/2$ is given by

$$\Delta a = c_\gamma^2\, c_w^2\, \frac{\alpha}{\pi}\, I(r), \tag{2.16}$$

where

$$I(r) = \frac{1}{2} - r + \frac{r}{2}(r-2)\ln(r) - \frac{r}{2}\frac{(r^2-4r+2)}{\sqrt{r^2-4r}} \ln \frac{r+\sqrt{r^2-4r}}{r-\sqrt{r^2-4r}}, \tag{2.17}$$

and $r = m_B^2/m_{\text{lepton}}^2$. Since r is large, we use the asymptotic form $I(r) \approx 1/(3r)$. Then the limit on c_γ corresponding to a two standard deviation uncertainty in the anomalous magnetic moment is

$$c_\gamma(m_\mu) < 0.050 \left(\frac{m_B}{\text{GeV}}\right) \tag{2.18}$$

for the muon, and

$$c_\gamma(m_e) < 0.360 \left(\frac{m_B}{\text{GeV}}\right) \tag{2.19}$$

for the electron. Thus, the constraint on c_γ for $m_B > 20$ GeV is roughly two orders of magnitude weaker than the constraints that we obtained from R.

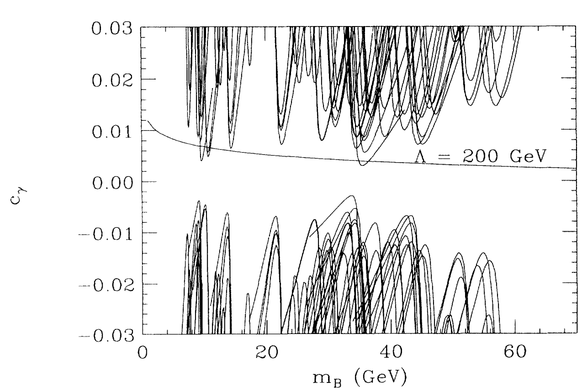

FIG. 3. Constraints on $c_\gamma(m_B)$ from the two standard deviations of the experimental uncertainty in R measured at various $\sqrt{s}$ as compiled by the Particle Data Group. The running of c_γ corresponding to $\Lambda = 200$ GeV is shown for comparison.

C. The scale Λ

It should now be clear that our original estimate of the sizes of $c_\gamma(\mu)$ and $c_Z(\mu)$ falls within the bounds that we have obtained from consideration of precision electroweak measurements. Recall that the estimate that we presented at the beginning of this section was for $c_\gamma(\Lambda) = c_Z(\Lambda) = 0$ at $\Lambda = m_{\text{top}}$. We will now determine how high we can push up Λ before we have unambiguous conflict with the experimental constraints. Using the approximate bound $c_Z(m_Z) < 0.02$, and assuming $m_{\text{top}} \approx 175$ GeV, we find

$$\Lambda < 1.3 \text{ TeV} \tag{2.20}$$

from Eqs. (2.3) and (2.4) with $\alpha_B = 0.1$. We can place comparable bounds on Λ from the constraints on c_γ, but the precise result depends crucially on the choice for m_B, as one can see from Fig. 3. What is interesting about Eq. (2.20) is that it implies that the scale of new physics lies at relatively low energies, just above the electroweak scale.

III. MODELS WITH NATURALLY SMALL KINETIC MIXING

In this section, we present a simple model with a gauged baryon number that naturally satisfies the boundary condition $c_\gamma(\Lambda) = c_Z(\Lambda) = 0$ with $\Lambda < 1.3$ TeV. There are two ingredients that are of central importance in the class of models that have small kinetic mixing below the electroweak scale. (1) In the full theory, at high energies, the kinetic mixing term is forbidden by gauge invariance. This is the case, for example, if we embed one of the U(1)s in a larger non-Abelian group. The mixing term remains vanishing down to the scale at which the gauge symmetry breaks to $G \times \mathrm{U}(1)_B \times \mathrm{U}(1)_Y$, where G contains the remaining gauge structure of the theory. (2) Beneath this symmetry-breaking scale, the one-loop diagram that connects the γ_B to the hypercharge gauge boson vanishes identically, so that c_γ and c_Z do not run. This places a constraint on the particle content beneath the symmetry-breaking scale,

$$\mathrm{Tr}\,(BY) = 0, \tag{3.1}$$

where B and Y are baryon number and hypercharge matrices, in the basis spanning the entire particle content of the theory. When we go to lower energies and the heaviest particle that contributes to Eq. (3.1) is integrated out, we will generate mixing through radiative corrections, in the way described quantitatively in Sec. II.

In what follows, we will present one example of a model with gauged baryon number that is "realistic" in the following sense: (i) the kinetic mixing is naturally small below the electroweak scale, (ii) there is a natural mechanism for generating the cosmic baryon asymmetry, and (iii) proton decay is forbidden (up to the usual nonperturbative effects) even though $\mathrm{U}(1)_B$ is spontaneously broken. It is not our goal to study every aspect of the phenomenology of this particular model, but rather to demonstrate by example that it is possible to construct models with the features (i), (ii), and (iii). In addition, we show that there are new fermions in the models of interest that appear in chiral representations of $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$; these fermions develop electroweak scale masses $\gtrsim m_{\mathrm{top}}$, which ensures that the scale Λ is not too far above the electroweak scale. Moreover, this implies that detection of the new fermions in this class of models is likely at an upgraded Tevatron or at the CERN Large Hadron Collider (LHC).

A. A model

The gauge structure of the model is

$$\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y \times \mathrm{SU}(4)_H,$$

where $\mathrm{SU}(4)_H$ is a horizontal symmetry. In addition to the ordinary three families of the standard model, f^i $(i = 1,2,3)$, we assume there is a fourth family F; the horizontal symmetry acts only on the quarks in the four families, which together transform as a **4** under the $\mathrm{SU}(4)_H$. The $\mathrm{U}(1)_B$ gauge group is embedded into $\mathrm{SU}(4)_H$ as

$$B = \begin{pmatrix} 1/3 & & & \\ & 1/3 & & \\ & & 1/3 & \\ & & & -1 \end{pmatrix}. \tag{3.2}$$

While $\mathrm{SU}(4)_H$ is broken at some high scale M_H, as we discuss below, the $\mathrm{U}(1)_B$ subgroup remains unbroken down to the electroweak scale. It is easy to verify that the particle content and quantum number assignments render the model anomaly-free.

We will also assume that there are right-handed neutrinos in each of the families. The right-handed neutrinos in the ordinary families acquire Majorana masses at a high scale M_N, while the one in the fourth family does not.[1] The fourth family neutrino develops an electroweak scale Dirac mass, so that the constraint from the invisible decay width of the Z is evaded. The Majorana masses for the right-handed neutrinos in the first three families are crucial to the baryogenesis scenario that we present in the next subsection. An interesting choice for the Majorana mass scale is $M_N \sim 10^{10}$–10^{12} GeV, which is consistent with the Mikheyev-Smirnov-Wolfenstein (MSW) solution to the solar neutrino problem, and the possibility that ν_τ is a hot dark matter particle [6].

The horizontal symmetry $\mathrm{SU}(4)_H$ is broken at a scale M_H down to $\mathrm{U}(1)_B$. One can imagine that the symmetry breaks in one step if a number of adjoint Higgs bosons, that we will generically call Φ, develop vacuum expectation values (VEV's) at the scale M_H. However, it is possible to generate a hierarchy of fermion Yukawa couplings if we break the symmetry sequentially, in the presence of additional vectorlike fermions [7]. The basic idea is as follows: We first introduce vectorlike fermions Ψ with mass M that transform as **4**s under $\mathrm{SU}(4)_H$, and assume that the electroweak Higgs boson H is a singlet under the horizontal symmetry. Suppose that the Ψ have Yukawa interactions such as $\mathcal{L} = \bar{q}_L \Phi \Psi + \overline{\Psi} H u_R + M \overline{\Psi} \Psi$. Then, when the vectorlike fermions are integrated out below the scale M, we obtain dimension-5 operators of the form

$$\frac{1}{M} \bar{q}_L H \Phi u_R \, . \tag{3.3}$$

Notice when the Φ's develop VEV's, operators such as the one in Eq. (3.3) generate Yukawa couplings in the low-energy theory. Now imagine that the horizontal symmetry is broken first to $\mathrm{SU}(3)_H \times \mathrm{U}(1)_B$ at the scale M_H, and then the $\mathrm{SU}(3)_H$ is broken sequentially at lower scales down to nothing. Then the dimension-5 operators that we have described can give rise to a hierarchical pattern of Yukawa couplings. A detailed analysis of the fermion mass matrix in models with horizontal symmetry breaking is beyond the scope of this paper, and we refer the interested reader to the literature [8].

The low-energy particle content of our model below both M_H and M_N is listed in Table I. Here, B refers to the gauge quantum number under $\mathrm{U}(1)_B$, while L is an effective nonanomalous global symmetry below M_N. $(B-L)_{\mathrm{extra}}$ is another nonanomalous global symmetry acting on the particles in the fourth family.

It is easy to see that the kinetic mixing remains vanish-

[1] This is natural if there is another global or local $\mathrm{SU}(4)_H$ acting on the leptons, and if lepton number embedded as $L = \mathrm{diag}(1,1,1,-3)$ in $\mathrm{SU}(4)_H$ is broken by an order parameter with $L = -2$ [or **10** under $\mathrm{SU}(4)_H$]. However, none of the conclusions in this paper depends on whether or not there exists a horizontal symmetry for the leptons.

TABLE I. Particle content below the horizontal symmetry breaking scale M_H and the right-handed neutrino masses M_N.

	Particle	$\mathrm{SU}(3)_C$	$\mathrm{SU}(2)_L$	$\mathrm{U}(1)_Y$	B	L	$(B-L)_{\mathrm{extra}}$
	q_L^i	3	2	1/6	1/3	0	0
Ordinary families	u_R^i	3	1	2/3	1/3	0	0
f_L^i, f_R^i	d_R^i	3	1	−1/3	1/3	0	0
$(i=1,2,3)$	l_L^i	1	2	−1/2	0	1	0
	e_R^i	1	1	−1	0	1	0
	Q_L	3	2	1/6	−1	0	−1
	U_R	3	1	2/3	−1	0	−1
Extra family	D_R	3	1	−1/3	−1	0	−1
F_L, F_R	L_L	1	2	−1/2	0	−3	+3
	E_R	1	1	−1	0	−3	+3
	N_R	1	1	0	0	−3	+3

ing down to the weak scale. Above M_H, the mixing is not allowed because $\mathrm{U}(1)_B$ is embedded into the non-Abelian group $\mathrm{SU}(4)_H$. This implies that the orthogonality condition (3.1) is satisfied by the particle content of the full theory. As we cross M_H, presumably all fields whose mass terms are allowed by the gauge symmetry decouple, but the particles listed in the table do not because they belong to chiral representations of the gauge group below M_H. One can easily check that the orthogonality between Y and B remains true below M_H as well given the particle content in Table I. The mixing term is only generated below the masses of the particles in the extra family (which we will refer to generically as m_F) which originate from electroweak symmetry breaking. Therefore, the mixing term remains vanishing down to the weak scale, i.e., $\Lambda = m_F \sim m_{\mathrm{top}}$, and the boundary condition discussed in the previous section is naturally achieved.

B. Baryogenesis and proton stability

It is natural to wonder how a cosmic baryon asymmetry can be generated in a model in which baryon number is a gauge symmetry at high energies. On the other hand, it is natural to worry about proton decay considering that the baryon number gauge symmetry is spontaneously broken at low energies. We address these two issues in this subsection. With regard to baryogenesis, we will show that the generation of a lepton number asymmetry from the decay of the right-handed Majorana neutrinos in our model can lead to a nonvanishing baryon number for particles from the ordinary three families, even though the total baryon number of the Universe remains zero. We describe this mechanism in some detail, as well as other relevant cosmological issues. Afterwards, we demonstrate that proton decay is forbidden in the model, apart from the electroweak nonperturbative effects.

The first step in baryogenesis is that a lepton asymmetry is generated from the CP-violating decays of the right-handed Majorana neutrinos. If we take into account the effect of the electroweak anomaly at a temperature at or above the electroweak phase transition [9], chemical equilibrium leads to nonvanishing lepton and baryon numbers in both the ordinary and extra families. Finally, the quarks in the extra family decay into those of the ordinary families, so that a cosmic over-density of fourth-generation particles is avoided. The first step is exactly the one proposed in Ref. [10] (see also [11] for the supersymmetric case.). The Yukawa interactions coupling the right-handed neutrinos to the lepton doublets violate CP in general; thus, the decay of the right-handed neutrinos can generate a net lepton asymmetry.

The analysis of chemical equilibrium including the electroweak anomaly effect is more complicated than in the minimal standard model [12]. Note that the total lepton number

$$L = (N_l + N_e) - 3(N_L + N_E + N_N) \tag{3.4}$$

is nonanomalous contrary to the minimal case, and is broken only by the small Majorana masses of the left-handed neutrinos generated by the seesaw mechanism [13]. In addition, there are the nonanomalous conserved quantum numbers

$$B = \frac{1}{3}(N_q + N_u + N_d) - (N_Q + N_U + N_D) \tag{3.5}$$

and

$$(B-L)_{\mathrm{extra}} = -(N_Q+N_U+N_D)+3(N_L+N_N+N_E), \tag{3.6}$$

where "extra" refers to fourth generation particles. The decay of the right-handed neutrinos generates only L, while both B and $(B-L)_{\mathrm{extra}}$ remain vanishing. Y also remains vanishing by gauge invariance. Since the number densities of the various species are proportional to their chemical potentials at the lowest order, we can derive nontrivial relations by considering the constraints imposed by chemical equilibrium.

The chemical equilibrium due to the Yukawa interactions implies

$$\mu_q = \mu_u + \mu_H = \mu_d - \mu_H, \tag{3.7}$$

$$\mu_l = \mu_e - \mu_H, \tag{3.8}$$
$$\mu_Q = \mu_U + \mu_H = \mu_D - \mu_H, \tag{3.9}$$
$$\mu_L = \mu_N + \mu_H = \mu_E - \mu_H, \tag{3.10}$$

while the electroweak anomaly effect requires

$$9\mu_q + 3\mu_l + 3\mu_Q + \mu_L = 0. \tag{3.11}$$

Here, μ_i refers to the chemical potential of a particle of species i and H is the standard electroweak Higgs boson. Then we find

$$B \propto \frac{1}{3}(18\mu_q + 9\mu_u + 9\mu_d) - (6\mu_Q + 3\mu_U + 3\mu_D) = 0, \tag{3.12}$$
$$(B-L)_{\text{extra}} \propto -(6\mu_Q + 3\mu_U + 3\mu_D) + 3(2\mu_L + \mu_N + \mu_E) = 0, \tag{3.13}$$
$$Y \propto \frac{1}{6}(18\mu_q + 6\mu_Q) + \frac{2}{3}(9\mu_u + 3\mu_U) - \frac{1}{3}(9\mu_d + 3\mu_D) - \frac{1}{2}(6\mu_l + 2\mu_L) - (3\mu_e + \mu_E) - 2\mu_H = 0, \tag{3.14}$$
$$L \propto (6\mu_l + 3\mu_e) - 3(2\mu_L + \mu_E + \mu_N) \neq 0. \tag{3.15}$$

Solving these constraints, we obtain

$$N_q + N_u + N_d = -\frac{108}{137}L, \tag{3.16}$$
$$N_l + N_e = \frac{101}{137}L, \tag{3.17}$$
$$N_Q + N_U + N_D = -\frac{36}{137}L, \tag{3.18}$$
$$N_L + N_E + N_N = -\frac{12}{137}L. \tag{3.19}$$

Thus, we see explicitly that the nonvanishing lepton number generated by the decay of the right-handed neutrinos will be partially converted to a nonvanishing baryon number for particles from the ordinary families, as well as nonvanishing baryon and lepton numbers for particles from the extra family.

One potential cosmological problem with this scenario is that the particles from the extra family could overclose the Universe. The constraints from primordial nucleosynthesis imply that baryons in the ordinary families must have a present cosmic density in the range $\Omega_b h_0^2 = 0.010$–0.15, where $h_0 = 0.4$–1 is the reduced Hubble constant [14]. On the other hand, the quarks and leptons in the extra family have also acquired an asymmetry that will remain until the present. Based on the predicted ratio of these asymmetries, the new contributions to the cosmic density are

$$\Omega_{Q,U,D} = \frac{m_F}{m_p}\Omega_b, \tag{3.20}$$
$$\Omega_{L,E,N} = \frac{1}{3}\frac{m_F}{m_p}\Omega_b, \tag{3.21}$$

where m_p is the mass of the proton. Even in the extreme case where $\Omega_b = 0.01$, the fourth generation particles would overclose the Universe when $m_F \gtrsim 100$ GeV. One might hope that these fourth generation particles could be candidates for cold dark matter. However, there are very strong observational constraints against dark matter that is strongly interacting [15], charged [16], or composed of Dirac neutrinos [17].

Fortunately, this problem can be avoided because baryon number is spontaneously broken, and we can find a way to make the fourth generation particles decay. Suppose that $U(1)_B$ is broken by an electroweak-singlet Higgs field χ with the following quantum numbers under[2] $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_B$:

$$\chi(\mathbf{1},\mathbf{1},\mathbf{1})_{+4/3}.$$

Then the following dimension-5 operators are the only ones for the quarks that are consistent with the gauge symmetry below M_H:

$$\mathcal{L}_5 = \frac{1}{M_V}(\bar{q}UH\chi + \bar{Q}uH\chi^* + \bar{q}DH^*\chi + \bar{Q}dH^*\chi^*) + \text{H.c.} \tag{3.22}$$

(Of course, there are similar operators involving the lepton fields.) One could imagine that these operators are generated by the exchange of a heavy vectorlike quark with mass M_V. As a consequence of Eq. (3.22), particles in the extra family can decay into ordinary particles. The decay rate is given by

$$\Gamma_F \sim \frac{1}{8\pi}\left(\frac{v\langle\chi\rangle}{M_V m_F}\right)^2 m_F, \tag{3.23}$$

where $v = 246$ GeV is the expectation value of the electroweak Higgs boson H, and $\langle\chi\rangle$ is the scale of baryon number symmetry breaking which we assume is around the electroweak scale. It is now clear that the particles in the extra family can decay well before nucleosynthesis as long as $M_V \lesssim 10^{15}$ GeV. Note that the decay of particles in the extra family gives an additional contribution to the cosmic baryon asymmetry of the ordinary particles that does not cancel out the original asymmetry that we obtained in Eq. (3.16).

The particles (especially quarks) in the fourth generation could be produced at the Tevatron or LHC. Their signatures depend on whether they leave the detector before or after they decay. If $M_V \lesssim 10^{11}$ GeV, they decay

[2] This charge assignment can be embedded into $SU(4)_H$ **15** representation, which allows the operator in Eq. (3.22).

inside the detector and leave a signature similar to that of the top quark. In this case, the fourth generation fermions must have masses larger than ~ 140 GeV [18]. If M_V is larger, they could be detected before they decay. A search for stable color-triplet quarks was carried out by the Collider Detector at Fermilab (CDF) [19], but the present constraint is rather weak (50–116 GeV ruled out at 95% C.L.).

An important point in this model is that proton decay is forbidden even though the baryon number is spontaneously broken. Any baryon-number-violating effects can be described in terms of effective operators with powers of the order parameter $\langle\chi\rangle$ which breaks $U(1)_B$. The general structure of such an operator is $O = q^k l^m \chi^{*n}$, where q is a quark field, l lepton, and k, m, n are integers. Since lepton fields carry integer charges and χ neutral, the power k has to be a multiple of 3, $k = 3l$ where l is another integer. Then the factor q^k carries baryon number l which has to be compensated by the baryon number of χ, $-(4/3)n$. Therefore n also has to be a multiple of 3, $n = 3p$, and the operator has the form

$$O = (q^{12}\chi^{*3})^p l^m, \tag{3.24}$$

which has dimension[3] $21p + (3/2)m$. Not only is this operator extremely suppressed by powers of a high mass scale, it also cannot contribute to proton decay because the quark field is raised to a power that is too large. Thus, there is no perturbative contribution to proton decay in this model. There could be a contribution from the electroweak instanton effect, but the decay rate due to the anomaly is known to be extremely tiny [20]. Therefore, nucleon decay is effectively forbidden in this model.

Finally, one might worry that γ_B exchange may lead to flavor-changing neutral current because it is coupled to a matrix diag$(1,1,1,-3)$ in the flavor space of the model. Mixing between the ordinary and extra families gives rise to off-diagonal coupling for the γ_B. However as seen above, the mixing is suppressed by a power of $\langle\chi\rangle/M_V$, and the off-diagonal coupling between different generations by a square of this suppression factor. All constraints from flavor-changing neutral currents are avoided when $M_V \gtrsim 100$ TeV. A similar lower bound applies to the mass of the horizontal gauge bosons in $SU(3)_H$.

IV. LEPTONIC SIGNALS

We have shown that there is a class of models with gauged baryon number in which the kinetic mixing between the hypercharge and baryon number gauge bosons is naturally small below the electroweak scale. Nevertheless, a small amount of mixing is not necessarily a bad thing, because it can provide us with a possible leptonic signature for our model. In this section we consider the new contribution to the Drell-Yan production of lepton pairs at hadron colliders. In particular, we show that the signal may be within the reach of an upgraded Tevatron.

The quantity of interest is $d\sigma/dM$, the differential cross section as a function of the lepton pair invariant mass. One can obtain the desired result from the conventional expression for the Drell-Yan differential cross section by making the substitutions

$$g_V^i \to g_V^i + \frac{2}{3} c_Z c_w s_w^2 \sqrt{\frac{\alpha_B}{\alpha}} \frac{\hat{s}}{\hat{s} - m_B^2 + i m_B \Gamma_B} \tag{4.1}$$

and

$$Q^i \to Q^i - \frac{1}{3} c_\gamma c_w \sqrt{\frac{\alpha_B}{\alpha}} \frac{\hat{s}}{\hat{s} - m_B^2 + i m_B \Gamma_B}, \tag{4.2}$$

where $\hat{s}$ is the parton center of mass energy squared. Here Q^i is the the quark charge in units of e, and $e g_V^i/(2 c_w s_w)$ is the vector coupling of the Z to a quark of flavor i, with $g_V^i = T_{3L} - 2Q^i s_w^2$.

Our results for $d\sigma/dM$ in a $p\bar{p}$ collision at $\sqrt{s} = 1.8$ TeV are shown in Fig. 4 for one lepton species, integrated over the rapidity interval $-1 < y < 1$, using the Eichten-Hinchliffe-Lane-Quigg (EHLQ) set II structure functions [21]. This range in rapidity was chosen to be consistent with the CDF detector coverage [22]. The solid line shows the conventional differential cross section (with $c_\gamma = c_Z = 0$), while the dotted lines give our results for $\alpha_B = 0.1$ and $c_\gamma(m_B) = c_Z(m_B) = 0.01$. For the values of m_B shown, the results do not depend strongly on the precise choice for c_Z. Around the γ_B mass there is a noticeable excess of events beyond the expected background. Because this excess is an interference effect, it depends linearly on c_γ. We show the excess in the total dielectron plus dimuon signal in a bin of size dM surrounding the γ_B mass in Table II, for $m_B = 30$, 40, and 50 GeV. The statistical significance of the signal

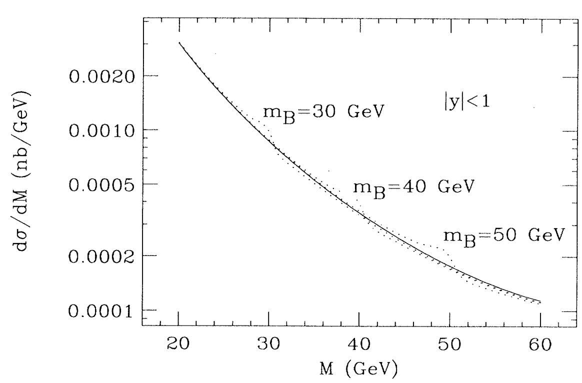

FIG. 4. Drell-Yan dilepton differential cross section as a function of the lepton pair invariant mass, integrated over the rapidity interval $|y| < 1$, for one lepton species. The dashed curves include the effect of γ_B exchange, assuming $\alpha_B = 0.1$ and $c_\gamma(m_B) = c_Z(m_B) = 0.01$, for $m_B = 30$, 40, and 50 GeV, respectively.

[3]This operator can be written in an explicitly $SU(2)_L \times U(1)_Y$ symmetric way: e.g., for $m = 0$, $p = 1$, $O = (q \cdot q)^4 d^4 \chi^{*3}$.

TABLE II. Excess dielectron plus dimuon production at the Tevatron, with $\alpha_B = 0.1$ and $c_\gamma(m_B) = c_Z(m_B) = 0.01$.

m_B	dM	Background	Excess	Statistical significance	
(GeV)	(GeV)	(fb)	(fb)	1 fb^{-1}	10 fb^{-1}
30	2	3468	320	5.4 σ	17.2 σ
40	4	2798	208	3.9 σ	12.4 σ
50	4	1422	112	3.0 σ	9.4 σ

assuming integrated luminosities of 1 fb^{-1} and 10 fb^{-1} is also shown. The largest excess at 1 fb^{-1}, is a 5.4 standard deviation effect for $m_B = 30$ GeV. However, with 10 fb^{-1} of integrated luminosity, even the excess at 50 GeV would be detectable at the 9.4 sigma level. This simple analysis is sufficient for a qualitative understanding of the signal we might expect to find at the Tevatron, with both the main injector, and a luminosity upgrade. We have not included the efficiency of the cuts and acceptance, but it is rather high even in a realistic analysis (93 % for e^+e^- and 82 % for $\mu^+\mu^-$ in CDF analysis [22]). A more exhaustive study, including the efficiency of the cuts and detector acceptance, as well as a comparison of the shape of the differential cross section to that expected in our model is required for a more accurate assessment of the discovery potential for this model at the Tevatron. It is interesting to note that even if the coupling α_B is smaller than 0.1 and the jet physics discussed in Ref. [1] is no longer of relevance, the mixing effect that we discuss here could still be significant enough to provide a clear signal for the model.

V. CONCLUSIONS

We have shown that there are models with a gauged baryon number in which kinetic mixing between the baryon number and hypercharge gauge bosons is naturally absent above the electroweak scale. Since the mixing is generated only through radiative corrections at lower energies, the resulting effective theory is consistent with precision electroweak measurements even when α_B is as large as 0.1, as we showed quantitatively in Sec. II. The exciting feature of the type of models that we proposed is that the baryon number gauge boson γ_B can be lighter than m_Z with a large gauge coupling, and yet be hidden in existing LEP and Tevatron data. This is the point that we emphasized in Ref. [1]. However, even if the coupling α_B is not large enough to produce an unambiguous hadronic signal, we have shown that the kinetic mixing term may give us another means for detecting the γ_B via its contribution to Drell-Yan dilepton production at hadron colliders. With both the main injector and a luminosity upgrade, this signal may eventually be within the reach of the Fermilab Tevatron.

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