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# Impact of extra particles on indirect $Z'$ limits

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We study the possibility of relaxing the indirect limits on extra neutral vector bosons by their interplay with additional new particles. They can be systematically weakened, even below present direct bounds at colliders, by the addition of more vector bosons and/or scalars designed for this purpose. Otherwise, they appear to be robust.

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## I. INTRODUCTION

In the coming years, the Large Hadron Collider (LHC) will explore energy scales up to a few TeV. New physics beyond the standard model (SM) at these scales could be unveiled, mainly in the form of new particles that give rise to resonances in the $s$ channel. One appealing possibility is the production of extra neutral vector bosons $Z'$[1] (for a recent review, see [1]). These particles appear in many extensions of the SM, such as grand unified theories (GUT), scenarios with strong electroweak symmetry breaking, theories in extra dimensions, and little Higgs models. $Z'$ bosons stand among the best candidates for an early discovery at the LHC. If coupled to quarks and leptons, they can be easily observed as dilepton peaks in the Drell-Yan process $q\bar{q} \to Z' \to \ell^+\ell^-(\ell = e, \mu)$. On the other hand, the existence of $Z'$ bosons with adequate couplings could explain the possible discrepancies with the SM predictions at large colliders, such as the anomalies recently found at the Tevatron: a $3.4\sigma$ discrepancy in the top forward-backward asymmetry for large invariant masses [2] (see [3]), or the $3.2\sigma$ excess in the $W + jj$ distribution around $M_{jj} \sim 150$ GeV [4] (see [5]).

However, other experiments have already placed stringent bounds on $Z'$ masses and couplings. They narrow significantly the parameter space available for LHC searches, especially in the early stages of the LHC at 7 TeV center of mass energy and low luminosity. These bounds are of two kinds: direct and indirect. Until recently, the best direct limits on $Z'$ bosons came from searches at the Tevatron. The same processes and couplings that are being studied at the LHC are relevant in this case. For sizes of couplings as in GUT (SM), the Tevatron typically puts a lower bound of $\sim$800 GeV (1 TeV) on the $Z'$ masses [6,7]. In this regard, it is noteworthy that, with only $\sim$40 pb$^{-1}$ of luminosity, the first LHC data allow one to derive limits comparable to those from the Tevatron [8,9]. Moreover, in some cases LHC bounds are already slightly stronger than the Tevatron ones. On the other hand, there are strong indirect limits from electroweak precision data (EWPD) obtained from measurements at the $Z$ pole (LEP and SLC), at low energies (experiments on parity violation and neutrino scattering), and above the $Z$ pole (LEP 2 and Tevatron) [10] (see also [11,12]).

The usually quoted $Z'$ limits depend on the implicit assumption that there are no other new particles that could change the analysis. However, in most models with extra neutral vector bosons, the $Z'$ is accompanied by other particles, such as additional neutral vectors, extra charged vector bosons, exotic fermions (sometimes required for anomaly cancellation), scalars (which may get a vacuum expectation value [VEV]), etc. In many instances, these additional particles do have an impact on the bounds. Direct limits are relaxed if the $Z'$ can decay into other particles beyond the SM, thus increasing its width. This happens, for instance, in some $Z'$ supersymmetric models [13]. Additional new particles contributing to EWPD will also modify the indirect bounds. Typically, the inclusion of more particles just makes the limits more stringent; in this case, the limits obtained from the analysis of a $Z'$ alone are valid as conservative limits. But it is also possible that their contributions cancel some effects of the $Z'$, in such a way that the limits are relaxed. In this paper we explore this possibility. By looking at the systematics of the cancellations and studying particular examples, we shall be able to judge the robustness of the standard $Z'$ limits.

The observable $Z'$ effects involving the SM matter fermions $\psi$ can be parametrized by the $Z'$ physical mass $M_{Z'}$, width $\Gamma_{Z'}$, mixing $s_{ZZ'}$ with the $Z$, and couplings $g^{\psi}$ to the $5 \times 3$ SM fermion multiplets. For simplicity, we shall assume family universality of the $Z'$ couplings (although similar arguments would apply to nonuniversal scenarios),

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[1]Neutral vector bosons can appear as components of different irreducible representations of $SU(2)_L \times U(1)_Y$. We will concentrate on $SU(2)_L \times U(1)_Y$ (and color) singlets. This is what is commonly called a $Z'$ boson.

 

yielding eight parameters. Particular models or classes of models impose relations on the couplings, thus reducing the number of independent parameters. For example, if there is only one Higgs doublet $\phi$ (as in the SM), and there are no additional particles light enough to affect $\Gamma_{Z'}$, then $\Gamma_{Z'}$ and $s_{ZZ'}$ can be computed in terms of the coupling to the Higgs, $g^{\phi}$, and the other parameters,[2] for a total of seven. For example, the mixing is given by

$$s_{ZZ'} \approx \frac{g^{\phi}\sqrt{g^2 + g'^2}}{2}\frac{v^2}{M_{Z'}^2}.$$

The fermion and Higgs couplings may also be related by additional assumptions, such as that the standard Yukawa couplings are allowed. Similarly, in specific GUT models the couplings to fermions are given by products of fixed charges times a coupling constant $g_{Z'}$, which is determined by the unification condition. The coupling to the scalar doublet or doublets may be fixed as well. In some cases this suffices to determine the mixing, while in others it depends on the ratios of unknown VEVs. In the latter case, the GUT $Z'$ model has two free parameters: $M_{Z'}$ and $s_{ZZ'}$.

EWPD are sensitive to the ratios $g^F/M_{Z'} \equiv G^F$, where $F$ represents any SM fermion or Higgs field. They also depend on the Higgs mass $M_H$ through the oblique radiative corrections. Including it, the model-independent fits with one $Z'$ and a single Higgs have basically eight free parameters,[3] whereas fits for GUT models with free mixing have three. The leading corrections to $Z$-pole observables depend linearly on the products $G^F G^{\phi}$. When the fermion couplings are fixed, this requires either large $M_{Z'}$ or small mixing $s_{ZZ'}$. On the other hand, the corrections to observables off the $Z$ pole depend also on the combinations $G^{\psi_1}G^{\psi_2}$, with $\psi_i$ any fermion multiplet. Therefore, these observables constrain the ratios $G^{\psi}$ even for vanishing mixing. For fixed fermion couplings, they put lower bounds on the $Z'$ mass. It is important to observe that the Tevatron and LHC searches and constraints depend on the fermion couplings and the mass, but not significantly on the mixing. Therefore, a vanishing mixing does not diminish the chances of observation at LHC.

The indirect limits on physics beyond the SM are in general rather stringent [14,15], in particular, on extra neutral vector bosons as already mentioned. This reflects two facts: the SM provides a quite good description of EWPD, and the largest deviations from the SM cannot be significantly accounted by such (simple) SM extensions. Hence, new particles, and in particular $Z'$s, are typically banished to high scales near a TeV. In the following we discuss more *complex* scenarios, where the new physics conspires to have little effect on EWPD, except maybe to accommodate a relatively heavy Higgs.[4]

In particular, we study the possibility of lowering the most stringent indirect limits on popular $Z'$ models by adding a second $Z'$ or/and extra scalars. In general, we can always lower those limits with specific additions designed for that purpose. Otherwise, the EWPD limits are robust. In Sec. II we introduce the effective dimension-six operators describing the $Z'$ contributions to EWPD observables. The effective Lagrangian approach is especially well-suited for the comparison of the contributions of any new heavy physics addition because it provides a basis of operators parametrizing any weakly coupled SM extension. Then, after discussing the different $Z'$ effects, we present the numerical analysis in Sec. III. We update the present indirect limits on popular $Z'$s with masses banished above a TeV, and characterize for each case the $Z'$ addition which may largely cancel the popular $Z'$ contributions to electroweak observables. In all cases we can lower those limits below the present Tevatron and LHC limits with properly chosen $Z'$s and scalars. For comparison, we also study the case of the minimal models discussed in [12], where the $Z$-$Z'$ mixing is not a free parameter. The corresponding constraints on the mixing can be somewhat relaxed by adding extra neutral and charged vector singlets. Our conclusions are collected in Sec. IV.

## II. EVADING ELECTROWEAK CONSTRAINTS

We want to examine which kind of additional new physics can relax the limits on $Z'$ bosons from EWPD. This is not straightforward, because the new particles that can neutralize the corrections of the $Z'$ to certain electroweak observables may simultaneously increase the discrepancy in others. A convenient way of analyzing the collective effect of several different particles is through their contribution to the gauge-invariant effective Lagrangian that describes arbitrary extensions of the SM at energies below the masses of the new particles. This procedure is more efficient than examining each of the many observables.

To analyze current EWPD it is sufficient to include dimension-four and dimension-six operators: $\mathcal{L}_{\rm eff} \approx \mathcal{L}_{\rm SM} + \mathcal{L}_6$. Here, $\mathcal{L}_{\rm SM}$ is the SM Lagrangian and

$$\mathcal{L}_6 = \frac{1}{\Lambda^2}\sum_i \alpha_i \mathcal{O}_i, \tag{1}$$

where $\mathcal{O}_i$ are gauge-invariant dimension-six operators, $\Lambda$ a scale of the order of the mass of the lightest extra particle, and $\alpha_i$ dimensionless coefficients. We will use the

[2]The Higgs mass $M_H$ affects the partial width for $Z' \to ZH$, but this effect is small for $M_{Z'} \gg M_H$. The effect of $M_H$ on the radiative corrections to EWPD are more important, as discussed below.

[3]Because the uncertainties for the other SM parameters are small, we fix them to their values at the SM minimum.

[4]As emphasized in Ref. [10] the main corrections to electroweak observables due to a relatively heavy Higgs can be canceled out by the addition of new vector bosons. Thus, as shown in Fig. 9 of that reference an extra vector boson triplet $\mathcal{W}^1$ or an adapted singlet $\mathcal{B}$ do balance the heavy Higgs contribution.

complete set of operators in [16] (see also [17]). In a given extension of the SM, the coefficients $\alpha_i$ can be written in terms of the couplings and ratios of masses of the new particles. The operators that contribute to EWPD can be classified into three groups.

(i) Oblique operators, which modify the $Z$ and $W$ propagators.
(ii) Operators made of scalars, gauge vector bosons (or derivatives), and fermion fields (scalar-vector-fermion [SVF] operators). They contribute to the trilinear couplings of two fermions and the $Z$ and $W$ bosons.
(iii) Four-fermion operators. Only the operators with at least two leptons contribute to EWPD.

At the $Z$ pole, only the oblique and SVF operators give significant contributions. The oblique operators also change the SM prediction for $M_W$. The high precision of these measurements strongly constrains the values of the oblique and SVF operator coefficients. Therefore, in practice only the four-fermion operators can give important contributions to observables off the $Z$ pole.

A sufficient condition for the new physics to be invisible to EWPD is that $\alpha_i$ vanishes (or is small enough) for all the operators contributing to electroweak observables. This condition is in practice also necessary, since different combinations of operators appear in the different observables. So, our strategy is simply to look for cancellations that make the relevant $\alpha_i$ vanish:

$$\alpha_i = \alpha_i^{Z'} + \alpha_i^{\text{extra}} = 0. \tag{2}$$

To start with, we need to know which of the operators that contribute to EWPD are induced by the $Z'$, and the values of their coefficients $\alpha_i^{Z'}$ in terms of the $Z'$ parameters. The following operators contributing to EWPD are generated in a generic $Z'$ model at tree level (see [10][5]):

(i) One oblique operator $\mathcal{O}_\phi^{(3)} = (\phi^\dagger D_\mu \phi)((D^\mu \phi)^\dagger \phi)$. After electroweak symmetry breaking this operator, induced by the $Z$-$Z'$ mixing, contributes to the $\rho$ parameter. Its coefficient is proportional to the $T$ parameter. No other oblique parameter appears in our operator basis at this order.
(ii) Five SVF operators, $\mathcal{O}_{\phi\psi}^{(1)} = (\phi^\dagger i D_\mu \phi)(\bar{\psi}\gamma^\mu \psi)$.
(iii) Nine four-fermion operators, $\mathcal{O}_{l\psi}^{(1)} = \frac{1}{1+\delta_{\psi l}} \times (\bar{l}_L \gamma^\mu l_L)(\bar{\psi}_L \gamma_\mu \psi_L)$, $\mathcal{O}_{l\psi} = (\bar{l}_L \psi_R)(\bar{\psi}_R l_L)$, $\mathcal{O}_{e\psi} = \frac{1}{1+\delta_{\psi e}} (\bar{e}_R \gamma^\mu e_R)(\bar{\psi}_R \gamma_\mu \psi_R)$, and $\mathcal{O}_{qe} = (\bar{q}_L e_R) \times (\bar{e}_R q_L)$.

Our notation is as follows: $\psi$ stands for any fermion multiplet; $\psi_{L,R}$ denote the left-handed (LH) doublets and right-handed (RH) singlets, respectively; $l_L$, $q_L$ represent the lepton and quark LH doublets; and $e_R$, $u_R$, $d_R$ the RH lepton and quark singlets. The SVF and four-fermion operators have two and four flavor indices, respectively, which we have not displayed. We have obtained the coefficients of these operators in [10]. The explicit expressions are collected in Appendix B of that reference (the $Z'$ is called $\mathcal{B}$ there). All the coefficients induced by a single $Z'$ have the factorized form

$$\alpha_{FF'}^{Z'} \propto g^F g^{F'}. \tag{3}$$

Because, in the universal scenario, there are only six different $Z'$ couplings $g^F$, there are many relations between the 15 operator coefficients.

Let us now examine which types of new particles can give the right contributions $\alpha_i^{\text{extra}}$ to offset, at least partially, the $Z'$ coefficients. We discuss, in turn, the different types of operators.

### A. Oblique operator

The $Z'$ contribution to the coefficient of the oblique operator $\mathcal{O}_\phi^{(3)}$ is negative definite: $(\alpha_\phi^{(3)})^{Z'} = -2(g^\phi)^2 \leq 0$. The contribution to the $\rho$ parameter has opposite sign, so it is positive definite. The same would hold, clearly, for any number of heavy $Z'$ bosons. This effect can be compensated for in several ways.

First, we need not resort to additional new physics. The loop effects of a heavy Higgs (with respect to the ones for a light Higgs, which is preferred by EWPD within the SM) give a negative correction to $\rho$ and can counterbalance to a large extent the $Z'$ contribution. In fact, if the Higgs were found to be heavy, a $Z'$ extension would be clearly favored over the SM. This mechanism has been analyzed quantitatively in [10]. It should be noted that the heavy Higgs also induces universal SVF operators, contributing to the $S$ parameter, so the cancellation is not perfect.

Second, a vanishing $\alpha_\phi^{(3)}$ can be achieved if the $Z'$ is accompanied by a singlet vector boson with hypercharge $Y = 1$, such as the one that appears in left-right models. This field gives rise, upon electroweak symmetry breaking, to an extra charged vector boson. Its main effect in EWPD is a negative contribution to $\rho$, proportional to the square of its coupling to the scalar doublet. In fact, this mechanism is at work in any model with a $Z'$ and custodial symmetry. On the other hand, its couplings to RH quarks (and to RH leptons if the neutrinos are Dirac) are constrained by measurements of $K^0$-$\bar{K}^0$ mixing, $\beta$ decay, $\mu$ decay, and weak universality [19,20].

Third, the effects can be cancelled or eliminated if several scalars participate in electroweak symmetry breaking. There are two distinct effects. One is that the $Z$-$Z'$ mixing is given for large $M_{Z'}$ by

[5]One can use other operator bases to describe the integration of $Z'$s [18], which may be more convenient for other purposes. We choose to use the standard basis [16,17], which is well adapted to perform global fits.

$$s_{ZZ'} \sim -\frac{\sum_i t_{3i} g^{\phi_i} \sqrt{g^2 + g'^2} |\langle \phi_i \rangle|^2}{M_{Z'}^2} \tag{4}$$

for an arbitrary set of scalar fields $\phi_i$ with $\{t_i, t_{3i}\}$ their weak isospin and third component, respectively. This can be reduced or eliminated by cancellations between $\phi_i$ with opposite signs for $t_{3i} g^{\phi_i}$, which are often present, e.g., in GUT models. A second effect is that the $\rho$ parameter is modified at the dimension-four level:

$$\rho^{\rm SM} \to \rho^{\rm multiple\ vevs} = \frac{\sum_i (t_i^2 - t_{3i}^2 + t_i) |\langle \phi_i \rangle|^2}{\sum_i 2 t_{3i}^2 |\langle \phi_i \rangle|^2}. \tag{5}$$

The $Z'$ contribution can be canceled by adjusting the VEVs of scalars with $t_i \geq 1$ and the appropriate $t_{3i}$.

Finally, the simplest possibility is making the $Z'$ scalar coupling $g^\phi$ and thus the mixing small. As we have mentioned before, this has no consequences for collider searches.

### B. SVF operators

All the coefficients of SVF operators are proportional to the $Z$-$Z'$ mixing, so the simplest way to get rid of them is, again, to have a sufficiently small mixing, either by cancellations, as in (4), or by a small $g^\phi$. It is also possible to cancel a nonvanishing $Z'$ contribution with additional particles, as we discuss next.

Extra vectorlike leptons and quarks contribute exclusively to SVF operators via their mixing with the SM fermions [21,22]. Using the results of [23,24], it is easy to check that for any $Z'$ couplings there exist combinations of extra leptons and quarks, with adjusted Yukawa couplings, such that the five net SVF coefficients vanish. However, in generic cases the combinations are very contrived: several multiplets in different representations of the SM gauge group are required. Moreover, to avoid flavor-changing neutral currents a replica of the multiplets is necessary. This is similar to the discussion of scalars in Sec. II D.

A cleaner cancellation is achieved by adding other vector bosons that also mix with the $Z$ boson. According to the results in [10], only the extra vector bosons of vanishing hypercharge that are $SU(2)_L$ and color singlets can give contributions to the SVF operators generated by the $Z'$. Hence, we have to resort to additional $Z'$ bosons. The net contribution of $N$ $Z'$ bosons, including the original one, to the SVF coefficients will vanish if the following equations are satisfied:

$$\frac{\alpha_{\phi\psi}^{(1)}}{\Lambda^2} = -\sum_{n=1}^{N} G_n^\psi G_n^\phi = 0, \qquad (\psi = l, q, e, u, d), \tag{6}$$

with $G_n^F = g_n^F / M_n$ and $g_n^F, M_n$ the couplings and masses of the $n$th $Z'$. Already with $N = 2$, for any fixed couplings $g_1^\psi$, $g_1^\phi$ and fixed mass $M_1$, there are nontrivial solutions that make all these coefficients vanish: $G_2^\phi = \pm c_{\rm SVF} G_1^\phi$, $c_{\rm SVF} G_2^\psi = \mp G_1^\psi$, $\psi = l, q, e, u, d$, with $c_{\rm SVF}$ any real number. We call such a companion $Z'$ boson a *mirror* $Z'$.

Of course, the additional $Z'$ bosons increase the deviation in the $\rho$ parameter, but this can be taken care of as discussed above. Then, all the $Z$-pole observables will be blind to this pair of $Z'$ bosons. All these $Z'$ bosons also contribute to four-fermion observables, as we discuss in the next subsections.

### C. Indefinite-sign four-fermion operators

Scalar and vector bosons both contribute to four-fermion operators. To cancel most of the contributions of the $Z'$, the better suited fields are again additional neutral vector bosons, due to their chirality structure and quantum numbers. However, in the universal scenario the contribution of each $Z'$ to the four-fermion operators $\mathcal{O}_{ll}^{(1)}$ and $\mathcal{O}_{ee}$ is negative semidefinite, so the different contributions go in the same direction and cannot cancel. The same holds for vector bosons with other quantum numbers [10]. Then the question is whether canceling the coefficients of these two operators is possible by introducing scalar fields. We will postpone that analysis to the next subsection, and concentrate here on the seven four-fermion operators with coefficients of indefinite sign.

Let us consider again a set of $N$ $Z'$ bosons. The cancellations we are looking for are given, in this sector, by nontrivial solutions to the following system of equations:

$$\frac{\alpha_{lq}^{(1)}}{\Lambda^2} = -\sum_{n=1}^{N} G_n^l G_n^q = 0, \tag{7}$$

$$\frac{\alpha_{lu}}{\Lambda^2} = 2\sum_{n=1}^{N} G_n^l G_n^u = 0, \tag{8}$$

$$\frac{\alpha_{ld}}{\Lambda^2} = 2\sum_{n=1}^{N} G_n^l G_n^d = 0, \tag{9}$$

$$\frac{\alpha_{le}}{\Lambda^2} = 2\sum_{n=1}^{N} G_n^l G_n^e = 0, \tag{10}$$

$$\frac{\alpha_{eu}}{\Lambda^2} = -\sum_{n=1}^{N} G_n^e G_n^u = 0, \tag{11}$$

$$\frac{\alpha_{ed}}{\Lambda^2} = -\sum_{n=1}^{N} G_n^e G_n^d = 0, \tag{12}$$

$$\frac{\alpha_{qe}}{\Lambda^2} = 2\sum_{n=1}^{N} G_n^q G_n^e = 0. \tag{13}$$

For given $G_1^\psi$ of the initial $Z'$, the unknowns are the ratios $G_n^\psi$, $n \geq 2$. In the case of two $Z'$s, $N = 2$, there are

seven equations for five unknowns. In fact, Eqs. (11) and (12) are not independent for $N = 2$, and will be automatically satisfied if Eqs. (7)–(9) and (13) are. However, it is easy to see that Eqs. (7)–(11) have no real solution unless $G_1^l = 0$, $G_1^e = 0$, or $G_1^q = G_1^u = G_1^d = 0$. Nevertheless, the EWPD limits can be relaxed if the nonvanishing coefficient only enters observables that are less precisely measured. With $N = 3$, all seven coefficients of indefinite sign can be zero.

On the other hand, if the $Z$-$Z'$ mixings do not vanish, and assuming no other contributions to four-fermion or SVF operators, the system of $Z'$s would have to satisfy at the same time Eqs. (7)–(13) and (6). This requires at least four $Z'$s, including the initial one. In view of such a proliferation, it is important to keep in mind that each of the extra $Z'$s increases the size of the four-fermion operator coefficients of definite sign.

### D. Definite-sign four-fermion operators

We address now the operators with four LH leptons, $\mathcal{O}_{ll}^{(1)}$, and with four RH leptons, $\mathcal{O}_{ee}$. Assuming diagonal and universal couplings, the contribution of the $Z'$s to their coefficients is

$$(\alpha_{ll}^{(1)})_{ijkl} = -\sum_{n=1}^{N} (G_n^l)^2 \delta_{ij}\delta_{kl},$$
$$(\alpha_{ee})_{ijkl} = -\frac{1}{2}\sum_{n=1}^{N} (G_n^e)^2 (\delta_{ij}\delta_{kl} + \delta_{il}\delta_{kj}), \tag{14}$$

which is indeed negative semidefinite.[6] We have displayed in this case the flavor indices, as they will be important in the discussion below. These operators contribute to the lepton cross sections and asymmetries at LEP 2, and to purely leptonic low-energy observables, such as parity violation in Møller scattering and neutral current neutrino electron scattering. As shown in [10], these operators cannot be canceled by other vector bosons. However, we will see that scalar fields can do the job. In general, we need two types of scalars.

The coefficient $\alpha_{ee}$ can be canceled by the contribution of a *scalar singlet* $\varphi$ with hypercharge $Y = -2$.[7] This scalar can only have the following renormalizable interaction with the SM fermions:

$$\Delta\mathcal{L} = -\lambda_\varphi^{ij} \varphi^\dagger \overline{e_R^{i\,c}} e_R^j + \text{H.c.} \tag{15}$$

The Yukawa coupling matrix $\lambda_\varphi$ is symmetric. From the interaction above, we see that $\varphi$ has lepton number $L = 2$. Integrating it out, the operator $\mathcal{O}_{ee}$ is generated, with coefficient

$$(\alpha_{ee})_{ijkl} = \frac{(\lambda_\varphi^\dagger)^{ki}\lambda_\varphi^{jl}}{M_\varphi^2}. \tag{16}$$

The coefficient is symmetric under exchange of the first and third and/or second and fourth family indices, just as the coefficient generated by the $Z'$s. This is actually a symmetry of the operator itself. Moreover, the coefficient has the right sign required to cancel the effect of the $Z'$s to this operator. The coefficient $\alpha_{ll}^{(1)}$ can, in turn, be canceled by the contribution of a *scalar triplet* $\Delta$ with hypercharge $Y = 1$. These scalars are well known by their role as messengers in the seesaw mechanism of type II [25]. For our purposes neutrino masses can be neglected and we can assume lepton number conservation. Thus, assigning to $\Delta$ lepton number $L = -2$, its only possible renormalizable interaction with SM fermions is

$$\Delta\mathcal{L} = -\lambda_\Delta^{ij} \overline{l_L^{i\,c}} i\sigma_2 \Delta^a \sigma_a l_L^j + \text{H.c.} \tag{17}$$

The contribution to the coefficient of the operator with four LH leptons is [26]

$$(\alpha_{ll}^{(1)})_{ijkl} = 2\frac{(\lambda_\Delta^\dagger)^{ki}\lambda_\Delta^{jl}}{M_\Delta^2}. \tag{18}$$

Again, this can neutralize the corresponding $Z'$ contribution.

There is, however, an important difficulty in realizing these cancellations. Because of the different chiral structure of their couplings, the flavor indices of the scalars are crossed with respect to the ones of the vectors. For the relevant observables, the first two (or last two) indices correspond to the first family, i.e., to electrons. Fixing these two indices in the operator coefficient, for the vector contribution one is left with a symmetric matrix, which in the diagonal universal case is proportional to the identity. In the scalar contribution, on the other hand, the coefficient reduces to a rank-one matrix. So, the cancellation of the contribution of any number of universal $Z'$ bosons with only one singlet and/or triplet scalar always leaves nonvanishing off-diagonal scalar contributions. Removing them from the purely leptonic four-fermion operators requires at least three scalars of each type.

The scalars couplings, in general, lead to lepton flavor-violating processes, for which there are stringent constraints. For example, the couplings of the triplet must obey $|\lambda_\Delta^{e\mu(e\tau)}||\lambda_\Delta^{ee}|/M_\Delta^2 < 1.2 \times 10^{-5}(1.3 \times 10^{-2})$ TeV$^{-2}$ from $\mu^-(\tau^-) \to e^-e^+e^-$, and $|\lambda_\Delta^{e\mu}||\lambda_\Delta^{e\tau}|/M_\Delta^2 < 1.7 \times 10^{-2}$ TeV$^{-2}$ from $\tau^- \to e^-e^+\mu^-$ [27]. All these restrictions are satisfied if each scalar couples the electron to just one lepton family, i.e., $e$, $\mu$, or $\tau$, and they do not mix. So, for a perfect cancellation of the operators $\mathcal{O}_{ee}$ and $\mathcal{O}_{ll}^{(1)}$ without flavor violation we would need to introduce three singlet scalars, $\varphi_{e,\mu,\tau}$, with nonvanishing Yukawa couplings $\lambda_{\varphi_e}^{ee}$, $\lambda_{\varphi_\mu}^{e\mu} = \lambda_{\varphi_\mu}^{\mu e}$, and $\lambda_{\varphi_\tau}^{e\tau} = \lambda_{\varphi_\tau}^{\tau e}$, respectively,

[6] For general (nondiagonal and nonuniversal) couplings only the coefficients with $i = k$ and $j = l$ are negative semidefinite.

[7] Note that this kind of scalar field must not get a VEV.

TABLE I. SM fermion charges, $Q^{\psi}_{Z'} = G^{\psi}_{Z'} M_{Z'}/g_{Z'}$, with $g_{Z'} = \sqrt{5/3}g'$ ($g_{Z'} = \sqrt{g^2 + g'^2}$) for the popular (minimal) $Z'$ models discussed in Sec. III. For the left-right model $\alpha = 1.53$. The charges for the $\overline{Z'}$ models are functions of the corresponding $Z'$ charges and chosen for convenience. See the text for details. The charges for the $Z'_\eta$ model are also included for completeness.

| Model | $l_L$ | $q_L$ | $e_R$ | $u_R$ | $d_R$ |
|---|---|---|---|---|---|
| $Z'_\chi$ | $\frac{3}{2\sqrt{10}}$ | $-\frac{1}{2\sqrt{10}}$ | $\frac{1}{2\sqrt{10}}$ | $\frac{1}{2\sqrt{10}}$ | $-\frac{3}{2\sqrt{10}}$ |
| $\overline{Z'}_\chi$ | $-\frac{1}{10}Q^l_{Z'_\chi}$ | $5Q^q_{Z'_\chi}$ | 0 | $5Q^u_{Z'_\chi}$ | $5Q^d_{Z'_\chi}$ |
| $Z'_I$ | $-\frac{1}{2}$ | 0 | 0 | 0 | $\frac{1}{2}$ |
| $\overline{Z'}_I$ | $-\frac{1}{10}Q^l_{Z'_I}$ | 0 | 0 | 0 | $2Q^l_{Z'_\chi}$ |
| $Z'_S$ | $\frac{2}{\sqrt{15}}$ | $-\frac{1}{4\sqrt{15}}$ | $\frac{1}{4\sqrt{15}}$ | $\frac{1}{4\sqrt{15}}$ | $-\frac{2}{\sqrt{15}}$ |
| $\overline{Z'}_S$ | $-\frac{1}{10}Q^l_{Z'_S}$ | $5Q^q_{Z'_S}$ | 0 | $5Q^u_{Z'_S}$ | $5Q^d_{Z'_S}$ |
| $Z'_{\rm LR}$ | $\sqrt{\frac{3}{5}}\frac{1}{2\alpha}$ | $\sqrt{\frac{3}{5}}\frac{1}{6\alpha}$ | $\sqrt{\frac{3}{5}}(-\frac{\alpha}{2}+\frac{1}{2\alpha})$ | $\sqrt{\frac{3}{5}}(\frac{\alpha}{2}-\frac{1}{6\alpha})$ | $-\sqrt{\frac{3}{5}}(\frac{\alpha}{2}+\frac{1}{6\alpha})$ |
| $\overline{Z'}_{\rm LR}$ | $-\frac{1}{10}Q^l_{Z'_{LR}}$ | $2Q^q_{Z'_{LR}}$ | $-\frac{1}{10}Q^e_{Z'_{LR}}$ | $2Q^u_{Z'_{LR}}$ | $2Q^d_{Z'_{LR}}$ |
| $Z'_R$ | 0 | 0 | $-\frac{1}{2}$ | $\frac{1}{2}$ | $-\frac{1}{2}$ |
| $\overline{Z'}_R$ | 0 | 0 | $-\frac{1}{10}Q^e_{Z'_R}$ | $2Q^u_{Z'_R}$ | $2Q^d_{Z'_R}$ |
| $Z'_{B-L}$ | $-\frac{1}{2}$ | $\frac{1}{6}$ | $-\frac{1}{2}$ | $\frac{1}{6}$ | $\frac{1}{6}$ |
| $Z'_\eta$ | $\frac{1}{2\sqrt{15}}$ | $-\frac{1}{\sqrt{15}}$ | $\frac{1}{\sqrt{15}}$ | $\frac{1}{\sqrt{15}}$ | $-\frac{1}{2\sqrt{15}}$ |
| $Z'_{\rm min}$ | $-\frac{1}{2}g_Y - g_{B-L}$ | $\frac{1}{6}g_Y + \frac{1}{3}g_{B-L}$ | $-g_Y - g_{B-L}$ | $\frac{2}{3}g_Y - \frac{1}{3}g_{B-L}$ | $-\frac{1}{3}g_Y + \frac{1}{3}g_{B-L}$ |

and three triplet scalars, $\Delta_{e,\mu,\tau}$, with nonvanishing Yukawa couplings satisfying analogous relations.

However, in the triplet case there is an additional problem: the same coupling $\lambda^{e\mu}_{\Delta}$ that enter the LEP 2 $e^+e^- \to \mu^+\mu^-$ observables also contributes to $\mu$ decay, and hence, indirectly, to other observables such as Cabibbo-Kobayashi-Maskawa universality. In the absence of further new physics affecting $\mu$ decay [28], this coupling is strongly constrained. This prevents the required cancellation of the $Z'$ effects on LEP 2 $e^+e^- \to \mu^+\mu^-$ data.

## III. RELAXING ELECTROWEAK LIMITS ON $Z'$ MODELS

So far we have determined what kind of additional new physics is required to cancel the effects of a $Z'$ boson in EWPD. We have seen that an optimal cancellation requires introducing many different particles with specific couplings. Such a complicated scenario is, of course, highly unnatural. However, in most cases it is possible to weaken the limits significantly using only a few extra particles, for instance, adding an appropriate second $Z'$, or scalar singlets and triplets, or both. In this section we give several examples with these two types of additions. We consider a few popular $Z'$ models and show numerically how the limits are relaxed with the help of these extra particles. In particular, we have chosen those examples where current limits exclude the corresponding $Z'$ masses below 1 TeV. This includes the $Z'_\chi$, $Z'_I$, and $Z'_S$ models, inspired in $E_6$ GUTs, the $Z'_{LR}$ from left-right models, the $Z'_R$, whose charges are given by the third component of $SU(2)_R$, and the $Z'_{B-L}$ model. The explicit fermionic charges, $Q^{\psi}_{Z'}$, for these models can be found in Table I, where we have also included the $Z'_\eta$ charges for later convenience. All these charges are related to the ratios $G^{\psi}_{Z'}$ entering in the electroweak fit by $Q^{\psi}_{Z'} \equiv G^{\psi}_{Z'} M_{Z'}/g_{Z'}$, with $g_{Z'} = \sqrt{5/3}g' \approx 0.46$. For a more detailed description of these models and their origin see [1]. Finally, we also discuss the effects of extra particles in the class of minimal $Z'$ models discussed in [12]. For these models, which we will denote by $Z'_{\rm min}$, the charges are normalized with $g_{Z'} = \sqrt{g^2 + g'^2} \approx 0.74$ [see Eq. (20) below]. In all cases we assume family independent gauge couplings.

### A. Extra particle additions canceling $Z'$ effects

For the popular $Z'$ examples the electroweak bounds are typically derived assuming free mixing [10,11]. Thus, the limits on their masses are determined by the constraints from observables off the $Z$ pole, i.e., by their contributions to four-fermion operators.[8]

Following the discussion in Secs. II C and II D such contributions can be canceled by adding extra $Z'$s and extra scalar fields. Thus, the addition of such particles should allow for weaker bounds in the selected examples. For each of the above mentioned $Z'$s we will consider a custom counterpart, denoted by $\overline{Z'}$, with charges chosen to attain (at least partially) such cancellations. The charges for these

[8]The corrections to $Z$-pole observables are proportional to the mixing of the $Z'$ with the $Z$ boson, $s_{ZZ'}$, and thus can be adjusted to zero.

TABLE II. 95% C.L. electroweak limits on the $Z'$ masses for some of the most popular models. We compare the limits obtained in the single $Z'$ scenario with those including extra $\overline{Z'}$ vectors and scalars. See text for details and Table I for the $Z'$ and $\overline{Z'}$ charges. For comparison, we also give in parentheses in the second column the most stringent limit from direct searches at the Tevatron and LHC [6,7,9].

| Model | 95% C.L. electroweak limits | | | | | | | | | | |
|---|---|---|---|---|---|---|---|---|---|---|---|
| | $Z'$ alone | | $Z' + \overline{Z'}$ | | $Z' + \varphi + \Delta$ | | | $Z' + \overline{Z'} + \varphi + \Delta$ | | | |
| | $M_{Z'}$ [GeV] | | $M_{Z'}$ [GeV] | $\frac{\lvert g_{\overline{Z'}}\rvert}{M_{\overline{Z'}}}$ [TeV$^{-1}$] | $M_{Z'}$ [GeV] | $\frac{\lvert\lambda_\varphi\rvert}{M_\varphi}$ [TeV$^{-1}$] | $\frac{\lvert\lambda_\Delta\rvert}{M_\Delta}$ [TeV$^{-1}$] | $M_{Z'}$ [GeV] | $\frac{\lvert g_{\overline{Z'}}\rvert}{M_{\overline{Z'}}}$ [TeV$^{-1}$] | $\frac{\lvert\lambda_\varphi\rvert}{M_\varphi}$ [TeV$^{-1}$] | $\frac{\lvert\lambda_\Delta\rvert}{M_\Delta}$ |
| $Z'_\chi$ | 1035 | (900) | 856 | 0.77 | 869 | 0.19 | 0.36 | 475 | 1.2 | 0.35 | [0.07, 0.46] |
| $Z'_I$ | 1144 | (842) | 878 | 1.3 | 993 | | 0.34 | 681 | 1.6 | | 0.38 |
| $Z'_S$ | 1162 | (871) | 895 | 0.79 | 994 | 0.17 | 0.35 | 434 | 1.4 | 0.38 | 0.49 |
| $Z'_{\rm LR}$ | 1206 | (959) | 605 | 1.9 | 1234 | 0.14 | 0.33 | 611 | 1.9 | 0.17 | 0.33 |
| $Z'_R$ | 1146 | (1006) | 697 | 1.9 | 1146 | 0.15 | | 451 | 2.5 | 0.36 | |
| $Z'_{B-L}$ | 1261 | (971) | | | 905 | 0.27 | 0.37 | | | | |

| Model | $Z'_I + Z'_\eta$ | | $Z'_I + Z'_\eta + \varphi + \Delta$ | | | |
|---|---|---|---|---|---|---|
| | $M_{Z'_I}$ [GeV] | $M_{Z'_\eta}$ [GeV] | $M_{Z'_I}$ [GeV] | $M_{Z'_\eta}$ [GeV] | $\frac{\lvert\lambda_\varphi\rvert}{M_\varphi}$ [TeV$^{-1}$] | $\frac{\lvert\lambda_\Delta\rvert}{M_\Delta}$ [TeV$^{-1}$] |
| $Z'_I, Z'_\eta$ | 1061 | 489 | 767 | 373 | 0.32 | 0.39 |

$\overline{Z'}$ models are also shown in Table I, below the ones for the corresponding $Z'$ example. Let us explain how we choose them in general. First, since we are interested in the contributions to observables off the $Z$ pole we can choose the effective Higgs charge of these particles to be zero. For two $Z'$s, we can exactly solve six of the Eqs. (7) to (13). Choosing $G_1^{q,u,d} = \mp c_{4F} G_2^{q,u,d}$ and $c_{4F} G_1^{l,e} = \pm G_2^{l,e}$, with $c_{4F}$ a real number, we can solve all but (10). It is convenient to choose $c_{4F}$ to be small, which typically ensures that the second $Z'$ would not have been observed in the dilepton searches at the Tevatron and LHC. A small $c_{4F}$ also minimizes the extra contributions to all four-lepton operators, and, in particular, avoids increasing too much the size of the operators with four LH or four RH leptons, which are definite negative. On the other hand, we cannot choose $c_{4F}$ to be too small. Large hadronic ratios would imply either large charges that could spoil perturbativity or too low masses, which would go against the effective Lagrangian expansion we use (and present limits on extra $Z'$s and other new particles). In practice, we increase the hadronic charges by a factor smaller than the one used to suppress the leptonic ones, and we restrict to order-one charges.

For the cancellation of the four-lepton operators with definite sign we will use the extra scalars $\varphi$ and $\Delta$ introduced in Sec. II D. To illustrate the possible effects without introducing (adjusting) too many extra parameters, we will restrict to a somewhat minimal scenario. We will consider three scalars singlets, $\varphi_\ell$, each of them coupling the electron to only one lepton family $\ell = e, \mu, \tau$, respectively, avoiding lepton flavor-violating constraints in this way. Moreover, we choose all their couplings in Eq. (15) to be equal $\lambda_{\varphi_\ell}^{e\ell} (= \lambda_{\varphi_\ell}^{\ell e}) \equiv \lambda_\varphi$, $\ell = e, \mu, \tau$. On the other hand, we will only add one $\Delta$ and with a nonzero coupling only to electrons, $\lambda_\Delta^{ee} \equiv \lambda_\Delta$, thus evading $\mu$ decay data constraints. It must be stressed that the choice of equal $\varphi$ scalar couplings does not allow for the exact cancellation of the four-lepton operators resulting from the $Z'$ integration, as they have a different coefficient relation by a factor of 2. Thus, $(\alpha_{ee})_{eeee} = 2(\alpha_{ee})_{ee\mu\mu}$ for a $Z'$ with diagonal and universal couplings, whereas $(\alpha_{ee})_{eeee} = (\alpha_{ee})_{ee\mu\mu}$ for our choice of $\varphi$ couplings [see Eqs. (14) and (16)]. If we had chosen the $\varphi$ scalar couplings to obtain such a cancellation, the numerical results (limits) below would not change much because the fits involve many other data and contributions, with the exception of the $Z'_R$ addition as we will comment when discussing this extended model. Obviously, allowing for many different arbitrary additions and couplings the $\chi^2$ can be relatively improved, but not its significance.

### B. Results and discussion

We have updated the fit to EWPD for each of the above mentioned popular $Z'$ examples [10], and performed new fits to the further extended models. For each case we compute the bounds on the $Z'$ mass and the new parameters controlling the interactions of the extra particles. The results are presented in Table II.[9] Prior to discussing

[9]In general we allow for arbitrary $Z$-$Z'$ mixing. In order for this to vanish in the models considered, however, an extended scalar sector allowing for cancellations between different unknown VEVs is required, as emphasized in Sec. II. In the fits below we take this mixing to be a free parameter, except for the minimal models which are discussed at the end, without introducing explicitly the corresponding extra scalars.

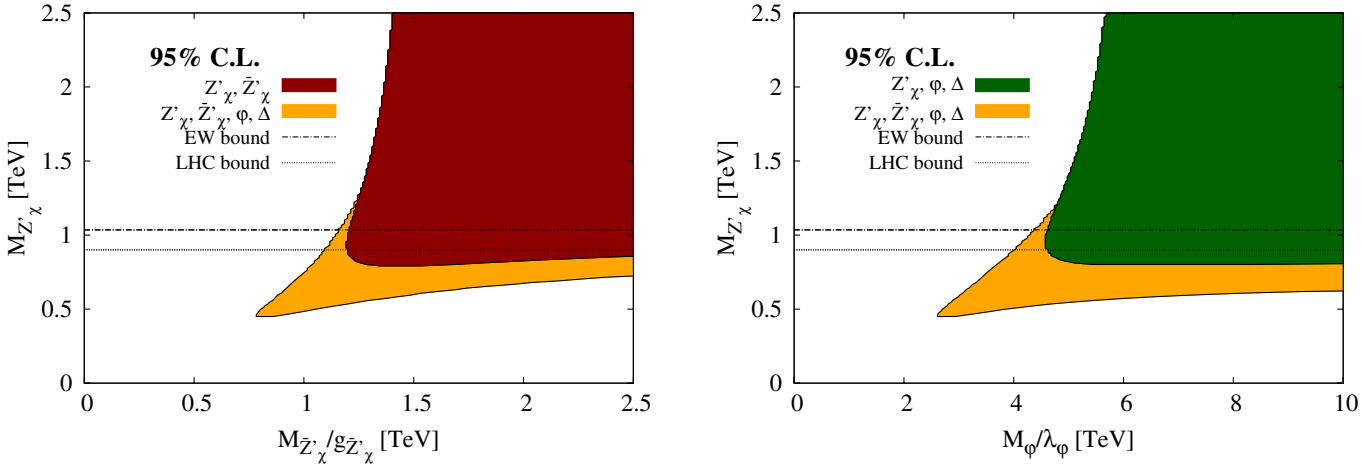


FIG. 1 (color online). Left: 95% C.L. confidence regions in the $M_{Z'_\chi} - (g_{\overline{Z'}_\chi}/M_{\overline{Z'}_\chi})^{-1}$ parameter space from a two-$Z'$ fit with and without including the scalars $\varphi$ and $\Delta$ (light [ocher] and dark [brown] solid regions, respectively). Right: the same in the $M_{Z'_\chi} - (\lambda_\varphi/M_\varphi)^{-1}$ plane from the fit to a $Z'_\chi$ plus the scalars $\varphi$, $\Delta$ (dark [green] solid region), and the fit including the $\overline{Z'}_\chi$ (light [ocher] solid region).

them, let us sketch a few details about the fits. We largely follow [10] and refer there for further details. There have been several updates of the experimental data considered in that reference. We use the updated values for the top mass [29], the SM Higgs direct searches limits from Tevatron [30], and the value for the five-quark contribution to the running of $\alpha_{\rm em}$ [31,32]. Other changes include the value of the $W$ width [33] and some updates in the low-energy data [20]. With these improvements the SM fit gives $M_H = 125$ GeV,[10] $m_t = 173.6 \pm 1.0$ GeV, $M_Z = 91.1876 \pm 0.0021$ GeV, $\alpha_S(M_Z^2) = 0.1184 \pm 0.0007$, and $\Delta\alpha^{(5)}_{\rm had}(M_Z^2) = 0.02751 \pm 0.00008$. At the minimum, $\chi^2_{\rm SM} = 158.5$ for a total of 207 degrees of freedom. As explained in the introduction, we fix all the SM parameters except the Higgs mass, which is left free, to these best SM fit values. Finally, the 95% C.L. limits (two-dimensional regions) presented in this section are obtained by requiring $\Delta\chi^2_{95\%} = 3.84(5.99)$ relative to the corresponding $\chi^2$ minimum for each case.

(i) $Z'_\chi$: The inclusion of the extra scalars suffices to pull the limit on $M_{Z'}$ for the $\chi$ model below the direct LHC bound of 900 GeV, even in the absence of any other $Z'$. In order to understand this, let us note that in this case the LEP 2 observables do not require an increase of the mass limit obtained from the $Z$-pole and low-energy data alone. On the contrary, they lower it (see Table 5 in [10]). This is because the effect of the $\chi$ model tends to increase the total $e^+e^- \to$ had cross section, which is $1.7\sigma$ above the SM. Thus, the addition of new scalars allows one to maintain this enhancement and compensate for the $Z'_\chi$ contributions to leptonic processes, which are in good agreement with the SM predictions. Furthermore, they also help to reduce the $1.3\sigma$ discrepancy with the electron weak charge extracted from Møller scattering. A similar limit is obtained if we consider the two $Z'$ scenario obtained by introducing the $\overline{Z'}_\chi$ in Table I. Note that in this particular case we have chosen not to couple the $\overline{Z'}_\chi$ to the RH leptons. Thus, we are not canceling any of the operators with RH quarks and leptons. This particular choice preserves large contributions to the hadronic cross section at LEP 2, while the cancellations prevent a significant discrepancy with the atomic parity violation data. Finally, when we include all the new particles, the limit on $M_{Z'_\chi}$ is lowered to around 475 GeV. The effect of adding different particles is illustrated in Fig. 1. On the left we show the 95% confidence region in the $M_{Z'_\chi} - M_{\overline{Z'}_\chi}/g_{\overline{Z'}_\chi}$ plane from a fit to the model with two $Z'$s alone (inner region), and with extra scalars in addition (outer region). On the right we show the corresponding regions in the $M_{Z'_\chi} - M_\varphi/\lambda_\varphi$ plane (in this case the inner one corresponds to the fit to $Z'_\chi$ alone plus the extra scalars). Apart from the scales, both figures look almost the same. Notice the significant correlation for low masses when we include all the particles at the same time. The correlation is less pronounced and the effect on the $M_{Z'_\chi}$ bound smaller for each separate addition. In particular, there is no

[10]Without including the direct limits the best-fit value for the Higgs mass still passes the barrier of 100 GeV, $M_H = 105^{+32}_{-26}$ GeV, getting closer to the LEP 2 exclusion bound of 114 GeV. This shift is due to the slightly larger value of the new top mass but mainly to the new determination of $\Delta\alpha^{(5)}_{\rm had}(M_Z^2)$ [32].

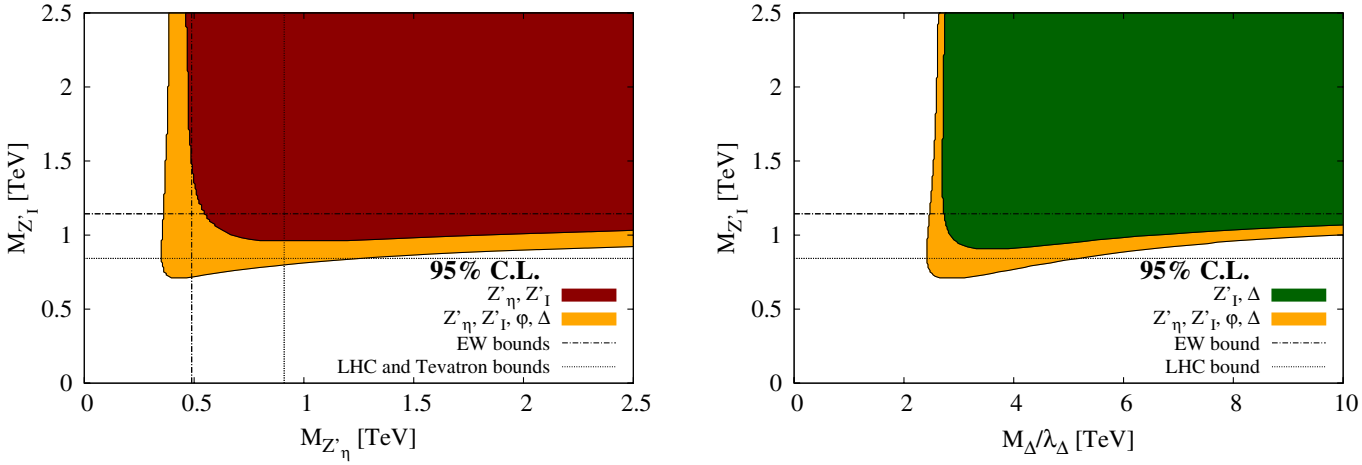


FIG. 2 (color online). Left: 95% C.L. confidence regions in the $M_{Z'_I}$-$M_{Z'_\eta}$ parameter space from the two-$Z'$ fit with and without including the scalars $\varphi$ and $\Delta$ (light [ocher] and dark [brown] solid regions, respectively). Right: the same in the $M_{Z'_I} - (\lambda_\Delta/M_\Delta)^{-1}$ plane from the fit to the $Z'_I$ plus the scalar $\Delta$ (dark [green] solid region) and from the fit also including $Z'_\eta$ and $\varphi$ (light [ocher] solid region).

appreciable correlation between the $M_{Z'_\chi}$ lower bound and $\varphi$, as observed in Fig. 1, right panel. In this case its significant reduction is mostly due to the triplet scalar $\Delta$ contribution.

(ii) $Z'_I$: As in the $Z'_\chi$ case, the limit on the $Z'_I$ mass can be slightly reduced adding new scalars (only the scalar triplet in this case, since $Z'_I$ does not couple to RH leptons). However, the new scalar is not enough to lower the EWPD limit below the LHC bound of 842 GeV, because the electroweak bound is more stringent in this case. The $Z'_I$ counterpart needs only to couple to LH leptons and RH $d$ quarks to attain a complete cancellation of all the four-fermion operators with no definite sign. The sole addition of the $\overline{Z'}_I$ in Table I lowers the limit around 100 GeV below that obtained with the scalar triplet. However, a complete cancellation of all four-fermion contributions is not possible when both $\overline{Z'}_I$ and $\Delta$ are included. This is because, even if we choose completely general couplings for $\Delta$, $\mu$ decay constraints on the electron-muon couplings prevent the cancellation of the $Z'$ contribution to the operator with two LH electrons and two LH muons. Still, the combined scenario in the right columns of Table II suffices to lower the electroweak limit significantly below the direct searches bound. We can also lower the $M_{Z'_I}$ limits with a second $Z'$ within $E_6$, $Z'_\eta$. (See Table II.) Actually, the charges of the $\eta$ and $I$ models are orthogonal within this group. When we combine these two $Z'$s with the (two) extra scalars, we can lower the limit on $M_{Z'_I}$ below the LHC bound of 842 GeV. This can be seen in Fig. 2. However, this limit occurs in correlation with a low $Z'_\eta$ mass and this is excluded below 910 GeV by Tevatron searches.[11] Taking this into account, the limit on $M_{Z'_I}$ is still slightly below the LHC bound.

(iii) $Z'_S$: The $Z'_S$ charges have a pattern rather similar to those of the $Z'_\chi$. Hence, we choose similar charges for its counterpart $\overline{Z'}_S$. Then, a similar discussion regarding the scalar additions and the combined scenario including all extra particles also applies, as can be seen in Table II.

(iv) $Z'_{LR}$: The limit on the $Z'_{LR}$ mass cannot be relaxed by introducing extra scalars. The reason is that for this model the LEP 2 constraints are dominated by the $e^+e^- \to$ had data and the effect of the $Z'$ is to reduce the total cross section relative to the SM, increasing the discrepancy with experiment.[12] We can ameliorate these restrictions, as well as those from low-energy data (in particular those from atomic parity violation experiments), by introducing a $\overline{Z'}_{\rm LR}$ with charges designed to cancel the left-right model contributions to operators with two leptons and two quarks. This addition alone suffices to lower the limit to half the electroweak bound in the single $Z'$ case. We find no improvement in this case when we also add extra scalars.

[11]We assume these limits still apply in the two-$Z'$ models, which is a good approximation if the resonances are narrow enough to be distinctively separated.

[12]Although it may seem surprising that in this case the 95% C.L. on the $Z'_{LR}$ mass in Table II is slightly higher when adding new scalars (parameters), this is so because this limit is relative to the corresponding new minimum. This is deeper due to the scalar contributions which do not decouple near this point, then redefining the probability distribution.

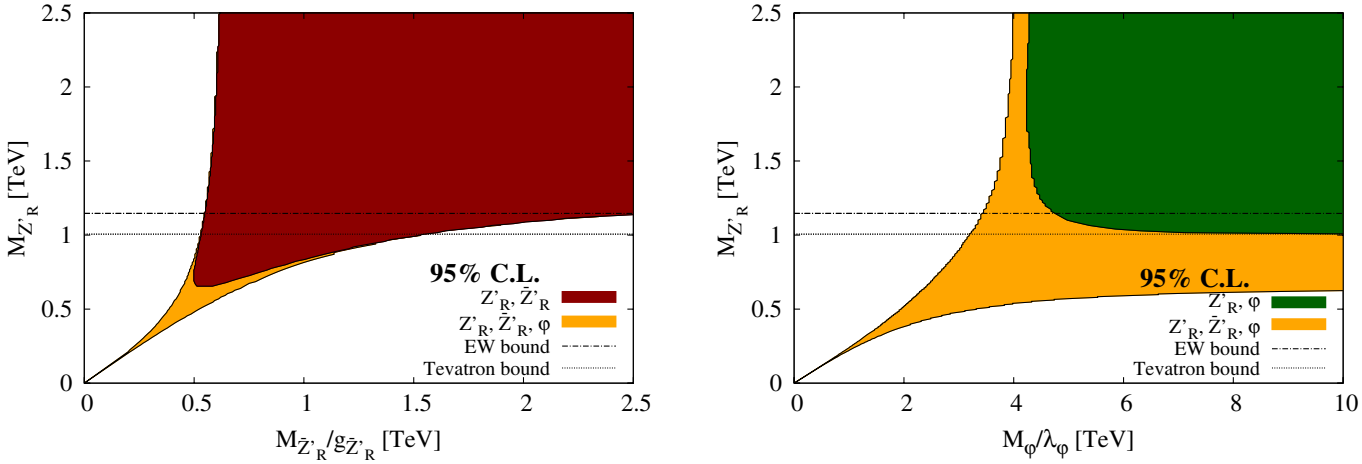


FIG. 3 (color online). Left: 95% C.L. confidence regions in the $M_{Z'_R} - (g_{\overline{Z}'_R}/M_{\overline{Z}'_R})^{-1}$ parameter space from a two-$Z'$ fit with and without including the scalars $\varphi$ (light [ocher] and dark [brown] solid regions, respectively). In this case, the scalar couplings are specifically chosen to attain a perfect cancellation of the purely leptonic $Z'_R$ effects. See text for details. Right: the same in the $M_{Z'_R} - (\lambda_\varphi/M_\varphi)^{-1}$ plane from the fit to a $Z'_R$ plus the scalars $\varphi$ (dark [green] solid region) and the fit including the $\overline{Z}'_R$ (light [ocher] solid region).

(v) $Z'_R$: Similar to the LR model, the limits on $Z'_R$ can be drastically reduced adding a $\overline{Z}'_R$. Also as in the LR case, the cancellation of the purely leptonic contributions by adding extra scalars alone leaves the limits intact. However, a significant improvement is possible when we combine the two additions. Since the $Z'_R$ only couples to RH fermions a complete cancellation of all four-fermion contributions would be possible with the addition of the second $\overline{Z}'_R$ and of scalar singlets $\varphi_\ell$ with couplings properly chosen. For our specific choice of scalar couplings, however, this cancellation is incomplete. Hence, we can find a 95% C.L. limit on $M_{Z'_R}$. As can be observed by comparing Eqs. (14) and (16), a perfect cancellation of the leptonic four-fermion operators requires that the scalar couplings satisfy the equality

$$\lambda^{ee}_{\varphi_e}/\sqrt{2} = \lambda^{e\mu}_{\varphi_\mu} = \lambda^{\mu e}_{\varphi_\mu} = \lambda^{e\tau}_{\varphi_\tau} = \lambda^{\tau e}_{\varphi_\tau}. \quad (19)$$

In such a case there is a flat direction in the parameter space, allowing for arbitrary $M_{Z'_R}$ values by adjusting the other extra parameters. This is illustrated in Fig. 3, which is analogous to Fig. 1 for $Z'_\chi$, but with the scalar coupling choice in (19). We must emphasize, however, that in the effective Lagrangian approach used here the fit only makes sense for $M_{Z'_R}$ above the maximum LEP 2 energies $\sim$209 GeV.

(vi) $Z'_{B-L}$: The limit on the $Z'_{B-L}$ mass is to a large extent determined by purely leptonic LEP 2 data. Thus, we do not find any $\overline{Z}'$ that can lower this limit. On the other hand, the addition of new scalars does allow for a $M_{Z'_{B-L}}$ limit around 350 GeV lower than in the single $Z'_{B-L}$ case.

The corresponding contours for the $Z'_S$, $Z'_{LR}$, $Z'_R$ (for our standard choice of $\varphi$ couplings), and $Z'_{B-L}$ are analogous to Figs. 1 and 2 for the $Z'_\chi$ and $Z'_I$.

Finally, we discuss the minimal $Z'$ models studied in [12]. Their charges are a linear combination of the hypercharge $Y$ and $B - L$. Thus, this case is fully characterized by the $Z'$ mass $M_{Z'_{\rm min}}$ and the two coupling constants $g_Y$ and $g_{B-L}$ defining its current. Following [12] we normalize these constants in such a way that the fermionic current coupling to the $Z'_{\rm min}$ is given (before mixing with the $Z$) by

$$J^\mu_{Z'_{\rm min}} \supset \sqrt{g^2 + g'^2} \sum_\psi [g_Y Y_\psi + g_{B-L}(B-L)_\psi] \bar{\psi} \gamma^\mu \psi. \quad (20)$$

Thus, the fit only constrains the ratios $g_Y/M_{Z'_{\rm min}}$ and $g_{B-L}/M_{Z'_{\rm min}}$. In Fig. 4 we depict the 95% confidence regions using two different parametrizations. On the left we draw the $M_{Z'_{\rm min}}/g_Y - M_{Z'_{\rm min}}/g_{B-L}$ plane to facilitate the comparison with previous cases, and on the right the $g_Y/M_{Z'_{\rm min}} - g_{B-L}/M_{Z'_{\rm min}}$ plane as done in [12]. In this class of models the relative sign between $g_Y$ and $g_{B-L}$ is physical. The limits are in general more stringent than the ones for the popular models above because by construction the $Z'$ mixing with the $Z$ boson is not a free parameter and is nonvanishing. Taking into account the different normalization used, which stands for a multiplicative factor $\sqrt{5/3}g'/\sqrt{g^2 + g'^2} \approx 0.6$, the lower limits in Fig. 4, left panel, are almost a factor $\sim$3 larger than those in the previous figures.

The results can be better visualized in Fig. 4, right panel, where the limits are plotted as a function of $g_Y/M_{Z'_{\rm min}}$ and $g_{B-L}/M_{Z'_{\rm min}}$. As in [12], we also draw the constraints from different data sets. In this class of models the addition of

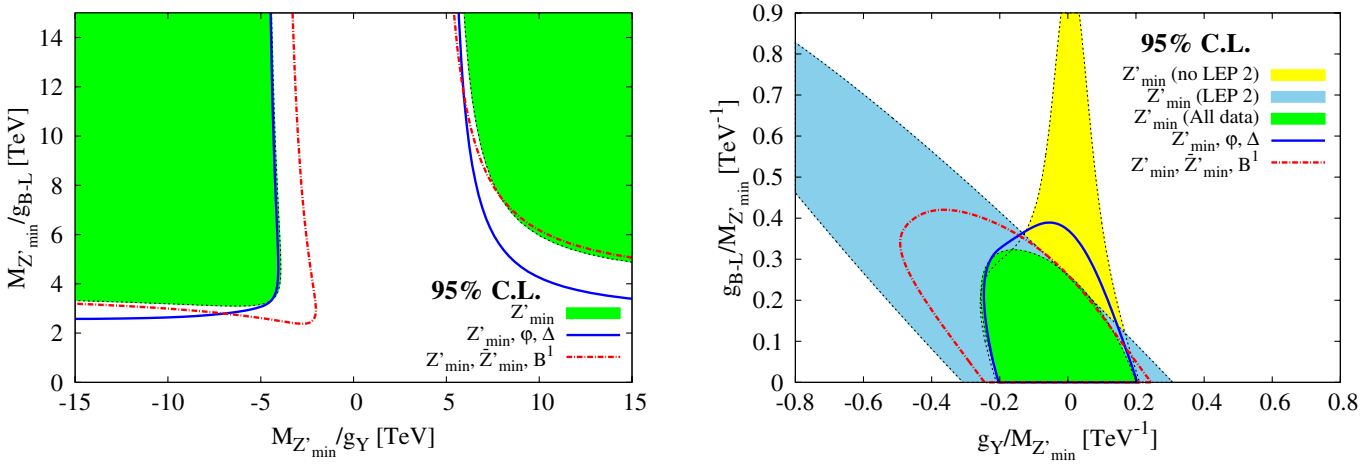


FIG. 4 (color online). Left: 95% C.L. confidence regions in the $M_{Z'_{\rm min}}/g_Y - M_{Z'_{\rm min}}/g_{B-L}$ parameter space for the minimal $Z'$ models (dark [green] solid region) and their extensions including also the scalars $\varphi$ and $\Delta$ (solid [blue] line) or a $\mathcal{B}^1$ together with a $\overline{Z'}$ with *mirror* charges (dot-dashed [red] line). Right: the same in the $g_Y/M_{Z'_{\rm min}} - g_{B-L}/M_{Z'_{\rm min}}$ plane. We also include the regions corresponding to the fit to EWPD without LEP 2 $e^+e^- \to \bar{f}f$ data (vertical [yellow] band) and to LEP 2 data only (diagonal [blue] band) for the $Z'_{\rm min}$ model alone.

new scalars has a limited effect. The corresponding 95% confidence region, which is delimited by the solid (blue) contours in Fig. 4, is not much larger than that of the minimal model alone. At any rate, due to the cancellation of the purely leptonic effects, the LEP 2 bounds can be somewhat relaxed and the allowed region is enlarged along the band determined by the $Z$-pole, $M_W$ and low-energy data (among others). When we include the new gauge bosons $\mathcal{B}^1$ and $\overline{Z'}_{\rm min}$,[13] the constraints from the $Z$-pole data and the observables sensitive to oblique effects can be significantly relaxed. Thus, the extended confidence region, delimited by the dot-dashed (red) contours in the figure, opens along the band allowed by LEP 2 data. The low-energy constraints, however, prevent one from obtaining a much larger region. On the other hand, (even) more contrived constructions including additional $Z'$s, in order to relax the remaining low-energy and LEP 2 hadronic constraints, may allow for larger regions.

## IV. CONCLUSIONS

New particles can manifest themselves as resonances at large colliders, or indirectly as deviations from the SM predictions in some observables. In this sense the Tevatron and LHC are complementary tools to EWPD for new physics searches, although none of them has provided significant evidence for physics beyond the SM yet. It is well known, however, that the quite good agreement of the SM predictions with EWPD implies that *simple* new physics is banished above the TeV scale, near the LHC reach. It is then important to investigate if more *complex* scenarios allow for small contributions to EWPD but relatively light new particles. This question is especially relevant to guide LHC searches.

In this paper we have addressed this question for extra neutral gauge bosons. We have discussed in turn several popular $Z'$ models based on $E_6$ and with EWPD mass limits above present direct bounds from the Tevatron and LHC. In particular, we have studied which additions of extra vector bosons and scalars can cancel their main contributions to EWPD. Using the effective Lagrangian approach, which is especially suited for comparing or combining different extensions of the SM, one can decide if there is a choice of couplings which may partially cancel the large contributions of any given extra gauge boson. We found that in all cases the EWPD bounds on the $Z'$ masses can be lowered below the present direct limits, although with specific additions designed for this purpose.[14] Otherwise, the EWPD limits appear to be robust.

We emphasize that the interest of the analysis presented here goes beyond the popular $Z'$ examples considered in Sec. III. Indeed, the methods and results in Sec. II are valid for arbitrary $Z'$ bosons with universal couplings, and the generalization to nonuniversal and flavor-changing couplings is straightforward.

If dilepton resonances are not found when more LHC data are available, the direct limits will eventually

[13] $\mathcal{B}^1$ is a fermiophobic singlet vector boson with hypercharge $Y = 1$, following the notation introduced in [10]; whereas $\overline{Z'}_{\rm min}$ is a $Z'$ with *mirror* minimal couplings, as described in Sec. II B. These additions allow for a complete cancellation of the $Z$-$Z'_{\rm min}$ mixing effects, as discussed in Secs. II A and II B. Of course, it is also possible to avoid these effects by allowing more general Higgs structures, as also described in those sections.

[14] We also require that the new particles have not been observed. In particular, that the custom $\overline{Z'}$ has a dilepton production cross section at large hadron colliders smaller than that of the $Z'$.

overcome those derived from EWPD for definite $Z'$s coupling to quarks and leptons, as long as the couplings are of the electroweak or GUT order. However, for large couplings the EWPD require large $Z'$ masses, which may be beyond the reach of the LHC (at least at 7 TeV). Indeed, $Z'$ production is suppressed at hadron colliders for large masses, due to the energy dependence of the parton distribution functions. This suppression is stronger than the $1/M_{Z'}$ scaling of EWPD limits. Thus, for large couplings the EWPD bounds may remain competitive.

In the opposite limit, leptophobic $Z'$ bosons [34] can be very light, since they evade Tevatron and LHC Drell-Yan bounds, and are not constrained by EWPD if they do not mix with the $Z$ boson. These $Z'$ bosons have been recently invoked [3,5] to account for Tevatron anomalies in the top forward-backward asymmetry [2] (other direct constraints apply in this case [35]) and $W + jj$ distribution [4]. Here we have shown that the EWPD constraints can also be evaded in models of leptophobic $Z'$ bosons with nonvanishing mixing. We note in passing that, because only the $W$ mass and $Z$-pole observables are relevant in this case, the effective Lagrangian approach can be accurate enough for a leptophobic $Z'$ with a mass $\sim$150 GeV.

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