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Interpretations of the ATLAS diboson resonances

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The ATLAS Collaboration has reported excesses in searches for resonant diboson production decaying into hadronic final states. This deviation from the standard model prediction may be a signature of an extra bosonic particle having a mass of around 2 TeV with a fairly narrow width, which implies the presence of a new perturbative theory at the TeV scale. In this paper, we study interpretations of the signal and its implication to physics beyond the standard model. We find that the resonance could be regarded as a leptophobic vector particle, which could explain a part of the observed excesses without conflict with the present constraints from other direct searches for heavy vector bosons at the LHC as well as the electroweak precision measurements.

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I. INTRODUCTION

Recently, the ATLAS Collaboration reported excesses in searches for massive resonances decaying into a pair of weak gauge bosons [1]. These excess events have been observed in the hadronic final states, i.e., the $pp \to V_1 V_2 \to 4j$ ($V_{1,2} = W^\pm$ or Z) channels. The weak gauge bosons from the resonance are highly boosted so that the hadronic decay products are reconstructed as two fat jets. Constructing the invariant mass of these two fat jets, it is possible to find a resonant peak for the intermediate state. The ATLAS collaboration has performed such an analysis by using 20.3 fb^{-1} data of the 8 TeV LHC running. Then, the excesses with narrow widths are observed around 2 TeV in the WZ, WW, and ZZ channels with local significance of 3.4, 2.6, and 2.9σ, respectively. Although we should wait for forthcoming ATLAS/CMS results of relevant searches to obtain a robust consequence about the observation, it should be worthwhile to consider possible interpretations of these anomalous events as evidence for new physics beyond the standard model (SM). In fact, the excesses are well fitted with resonances whose peaks are around 2 TeV and widths are less than about 100 GeV. Such narrow resonances may imply new weakly interacting particles, and then the underlying theories would be perturbative.[1] In this paper, we especially consider such a possibility to explain the excesses.

As mentioned above, the excesses reported by the ATLAS collaboration are in the WZ, WW, and ZZ channels. The tagging selections for each mode used in the analysis are, however, rather incomplete: about 20% of the events are shared by these channels. At the present stage, it may be hard to conclude that one resonance is responsible for the excesses in all the channels. There may be a possibility that one 2 TeV particle contributes to only one part of the channels and the peaks in the other channels are merely contamination due to the incomplete tagging selections. Taking this situation into account, in this paper, we do not limit ourselves to account for all of these excesses simultaneously, and consider the possibility that the new resonance appears in one channel. For each channel, the number of excess events could be accounted for if there is a 2 TeV resonance whose production cross section times decay branching ratio into gauge bosons is about 6 fb. We regard this as a reference value in what follows.

In order for a resonance to decay into two gauge bosons, it should be a bosonic state, namely, a particle with a spin zero or one under an assumption of the renormalizable theory. Let us first consider the spin-zero case. If such a particle is a singlet under the $\mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y$ gauge interactions, it couples to the electroweak gauge bosons and the SM fermions only through the mixing with the SM Higgs boson in the renormalizable potential. Therefore, its production cross section is suppressed by the mixing factor and in general too small to explain the anomalies. A possible way to enhance the production is to introduce new vectorlike colored particles. A singlet scalar field generically couples to these colored particles. Then the singlet is produced via the gluon fusion process according to the loop correction involving the vectorlike colored particles. If the masses of the vectorlike particles are above 1 TeV, such a scalar resonance with 2 TeV mass does not

[1]A possibility of strong dynamics is proposed and investigated in Ref. [2].

decay into these particles. It turns out, however, that $\mathcal{O}(10)$ fb production cross sections require $\mathcal{O}(10)$ extra colored particle pairs. Moreover, a large fraction of produced singlets decays into gluons, and thus gives only a negligible contribution to the diboson channels. Hence, a singlet scalar boson is inappropriate to explain the anomalies. An alternative possibility is to exploit $\mathrm{SU}(2)_L$ doublet scalars. These scalar fields may develop a finite vacuum expectation value (VEV) to directly couple to W and Z, and again mix with the SM Higgs field. To assure a large cross section to the diboson decay processes, there should be a sizable deviation from the SM limit. In addition, the deviation modifies the Higgs couplings, which are stringently constrained by the Higgs data at the LHC experiments. Within the constraint, both the production cross section and the branching ratios to the electroweak gauge bosons of a 2 TeV doublet scalar are found to be extremely small. A higher representation of $\mathrm{SU}(2)_L$ also suffers from its small production cross section since it does not couple to the SM quarks directly. For these reasons, we conclude that it is quite difficult to explain the required event rate with a new scalar particle, and thus we do not pursue this possibility in the following discussion.

Another candidate is a spin-one vector boson. Such a particle naturally appears if a high-energy theory contains additional gauge symmetry that is spontaneously broken at a certain scale above the electroweak scale. If the symmetry breaking occurs at the TeV scale, we expect the masses of the extra gauge bosons to be $\mathcal{O}(1)$ TeV. The gauge bosons are produced at the LHC if quarks are charged under the extra gauge symmetry. In this paper, we investigate this possibility. An important caveat here is that such a TeV-scale vector boson has been severely constrained by the LHC experiments. The strongest constraint is usually from the Drell-Yan processes [3,4]; for a 2 TeV vector boson, its production cross section times the branching ratio in the lepton final states should be much smaller than 1 fb. This bound makes it quite difficult to realize a sizable event rate for the diboson decay channel in most extensions of the SM with new gauge symmetries.

One promising setup to suppress the Drell-Yan processes is given by an $\mathrm{SU}(2)_L$ singlet heavy charged gauge boson with hypercharge ± 1, which we denote by W'. Such a W' is contained in some simple extensions of the SM, such as $\mathrm{SU}(2)_L \otimes \mathrm{SU}(2)_R \otimes \mathrm{U}(1)_{B-L}$ models [5]. This W' couples with right-handed quarks, as well as right-handed charged leptons and neutrinos. If right-handed neutrinos are rather heavy, W' is unable to decay to leptons and thus evade the Drell-Yan bounds. Since W' couples to the right-handed quarks, it is sufficiently produced at the LHC. After the electroweak symmetry breaking, the W' bosons mix with the electroweak gauge bosons, which allows W' to decay into W and Z. In Sec. II, we study whether W' can explain the observed excess. We find that it is difficult to realize large decay branch into WZ in a simple version of the W' model, and therefore the required event rate for the ATLAS diboson excess is not obtained once the limit from the electroweak precision measurements is taken into account. Even if this limit is avoided by canceling the W' contribution to the electroweak precision observables with other new physics effects, the resonance search in the channel consisting of a W boson and a Higgs boson (Wh) [6] severely constrains the WZ decay branch. Besides, W' generally predicts flavor changing gauge couplings [7], as the W boson does in the SM. We have to assume these couplings to be flavor diagonal to evade the strong bounds from flavor physics. This gives rise to additional complexity for a concrete model building in this direction.

An alternative way is to regard the resonance as a neutral massive gauge boson Z' which has no coupling to the SM leptons. This is the so-called leptophobic Z'. Such a leptophobic Z' may be realized in the grand unified theories (GUTs); if the rank of the GUT group is larger than four, it includes extra U(1) symmetries, and a certain linear combination of the U(1) charges could be leptophobic. Especially, a set of charge assignments inspired by the E_6 GUTs has been widely studied so far in the literature [8–12]. The Drell-Yan bounds on this class of models are then readily avoided because of the leptophobic nature. Again, Z' mixes with the Z boson after the electroweak symmetry breaking, and thus it has a decay mode into a pair of $W^{\pm}$. We study the decay properties of such a Z' using a simplified model in Sec. III to see whether it could explain the ATLAS diboson signal.

Finally, in Sec. IV, we conclude our discussion and give some future prospects for probing the scenarios in the future LHC experiments.

II. W' MODEL

To begin with, we consider a simplified model for W' to study whether it explains the ATLAS diboson signal or not. For recent works on phenomenological studies of W', see Ref. [13]. As mentioned in the introduction, we consider an $\mathrm{SU}(2)_L$ singlet vector boson with $+1$ hypercharge as a candidate for W'^{+}, since it effectively has no coupling to the SM leptons and thus avoids the severe Drell-Yan bounds. Such a vector boson may be attributed to a gauge boson of a non-Abelian gauge group orthogonal to $\mathrm{SU}(2)_L$, like $\mathrm{SU}(2)_R$. We may also take up an $\mathrm{SU}(2)_L$ triplet nonhypercharged vector boson, which, for instance, appears in the $\mathrm{SU}(2)_1 \otimes \mathrm{SU}(2)_2 \otimes \mathrm{U}(1)_Y$ type models [14]. In this case, however, couplings of the $\mathrm{SU}(2)_L$ triplet vector bosons to the SM charged leptons are generically allowed, and thus we need an additional mechanism to suppress these couplings to evade the Drell-Yan constraints. In this sense, the $\mathrm{SU}(2)_L$ singlet vector boson is more favored, and thus we focus on this candidate in our work.

Let us denote the massive $\mathrm{SU}(2)_L$ singlet vector boson by $\hat{W}'^{+}$. There are scalars charged under the additional SU(2) symmetry, and they develop nonzero VEVs to cause the SU(2) symmetry breaking. Then $\hat{W}'^{+}$ gains a TeV-scale mass. We assume that some of the scalars are charged under $\mathrm{SU}(2)_L$ as well and the finite mass mixing between $\hat{W}'^{+}$ and $\hat{W}^{+}$ in the SM is generated by their VEVs. The mix is described as

$$\begin{pmatrix} W^+ \\ W'^+ \end{pmatrix} = \begin{pmatrix} \cos\zeta & \sin\zeta \\ -\sin\zeta & \cos\zeta \end{pmatrix} \begin{pmatrix} \hat{W}^+ \\ \hat{W}'^+ \end{pmatrix}, \quad (1)$$

where W^+ and W'^+ are the mass eigenstates. We expect the mixing angle ζ is $\mathcal{O}(v^2/M_{W'}^2)$ where $M_{W'}$ is the mass of W' and $v \simeq 246$ GeV is the Higgs VEV. Then the partial decay width of W'^+ into W^+ and Z is given as follows:

$$\Gamma(W'^+ \to W^+ Z) \simeq \frac{\alpha_2 \sin^2 2\zeta}{192} \frac{M_{W'}^5}{M_W^4}, \quad (2)$$

where α_2 is the $\mathrm{SU}(2)_L$ gauge coupling and M_W is the mass of W boson. From this expression, we find that although the partial decay width is suppressed by the small mixing angle ζ, this suppression is compensated for by the enhancement factor of $(M_{W'}/M_W)^4$. This enhancement factor results from the high-energy behavior of the longitudinal mode of W'. Therefore, we expect a sizable decay branch for the W^+Z channel. The partial decay width gets increased as the mixing angle becomes large. The size of the mixing angle is, on the other hand, restricted by the electroweak precision measurements since it is induced by interactions which break the custodial symmetry—namely, the bound on the T parameter [15] constrains the mixing angle. The current limit on ζ is given by $|\zeta| \lesssim 5 \times 10^{-4}$ for $M_{W'} = 2$ TeV [16], which in turn gives an upper limit on $\Gamma(W' \to W^+ Z)$. We however note that the constraints may be evaded if there is another contribution to the T parameter which cancels the effects of the W-W' mixing. The actual realization of this possibility is model dependent, and we do not pursue it in this paper.

The equivalence theorem tells us that the final state gauge bosons in the $W'^+ \to W^+ Z$ channel could be regarded as Nambu-Goldstone bosons, since the longitudinal mode dominates the decay amplitude as we have just mentioned. Thus, the partial decay width of the channel is related to that of the decay to W^+ and the Higgs boson in the final state. In fact, we have

$$\Gamma(W'^+ \to W^+ h) \simeq \Gamma(W'^+ \to W^+ Z), \quad (3)$$

where h is the SM-like Higgs boson. Currently, the CMS Collaboration gives an upper bound on this decay mode [6] as $\sigma(pp \to W'^+) \times \mathrm{BR}(W'^+ \to W^+ h) \lesssim 7$ fb. Thus, through the above equation, this bound also implies $\sigma(pp \to W'^+) \times \mathrm{BR}(W'^+ \to W^+ Z) \lesssim 7$ fb, which somewhat conflicts with the ATLAS diboson anomaly. Since the above relation is a consequence of the equivalence theorem, this bound is robust and almost model independent. For this reason, a W' model (as well as a Z' model as we see in the next section) in general predicts a smaller number of signals in the diboson channel than that observed in Ref. [1], once we consider the limit on the Wh channel.

$\hat{W}'^{+}$ carries the $+1$ hypercharge, so that the $\mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y$ symmetry allows the following couplings the right-handed quarks:

$$\mathcal{L}_{W'ud} = \frac{g_{ud}}{\sqrt{2}} \bar{u}_i \hat{W}'^{+} P_R d_i + \mathrm{H.c.}, \quad (4)$$

where $P_{L/R} \equiv (1 \mp \gamma_5)/2$ and $i = 1, 2, 3$ denotes the generation index. Here, we assume the coupling constant g_{ud} is common to all of the generations and ignore flavor nondiagonal parts for brevity, which are in fact stringently constrained by the measurements of the flavor observables, such as the K^0-$\bar{K}^0$ mass difference. At the LHC, W' is produced via the interactions in Eq. (4). For W' with a mass of 2 TeV, the production cross section at the LHC with the center-of-mass energy $\sqrt{s} = 8$ TeV (LHC8) is evaluated as

$$\sigma(pp \to W'^{\pm}) \simeq 490 \times g_{ud}^2\ [\mathrm{fb}], \quad (5)$$

using MadGraph [17]. After the production, W' mainly decays into the WZ, Wh, or quark final states. The partial decay width for the final state containing a pair of u_i and $\bar{d}_i$ is given by

$$\Gamma(W'^+ \to u_i \bar{d}_i) = \frac{g_{ud}^2}{16\pi} M_{W'}, \quad (6)$$

where we neglect the quark masses for brevity. The branching fraction of this decay mode is severely constrained by the dijet resonance searches [18,19]. Following Ref. [18], we have $\sigma(pp \to W') \times \mathrm{BR}(W' \to q\bar{q}') \lesssim 100$ fb with the acceptance $\mathcal{A} \simeq 0.6$ being assumed. The ATLAS Collaboration gives a similar limit on the dijet channel [19]. Currently, the $W'^+ \to t\bar{b}$ decay mode is less constrained [20]: $\sigma(pp \to W'^+) \times \mathrm{BR}(W'^+ \to t\bar{b}) \lesssim 120$ fb.

Taking the above discussion into account, in Fig. 1, we show a contour plot for $\sigma(pp \to W') \times \mathrm{BR}(W' \to WZ)$ on the g_{ud}-ζ plain. The light-green shaded region is disfavored by the electroweak precision measurements: $|\zeta| \lesssim 5 \times 10^{-4}$. The dark-green and blue shaded regions are excluded by the limits from the $W' \to Wh$ [6] and the dijet [18,19] channels, respectively. This figure clearly shows that it is difficult to explain the observed diboson resonance with this simplified W' model once the electroweak precision bound on the W-W' mixing is taken into account. Even if this constraint is avoided by utilizing other new physics effects to cancel the W' contribution to the electroweak precision observables, the bound on the $W' \to Wh$ channel

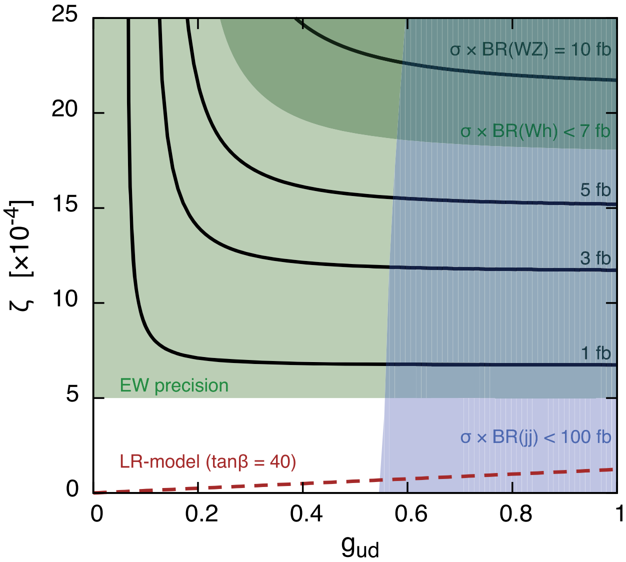

FIG. 1 (color online). $\sigma(pp \to W') \times \mathrm{BR}(W' \to WZ)$ (black solid lines). Here we set $M_{W'} = 2$ TeV. The light-green shaded region is disfavored by the electroweak precision measurements. Dark-green and blue shaded regions are excluded by the limits from the $W' \to Wh$ [6] and the dijet [18,19] channels, respectively. The brown dashed line represents the case of the $\mathrm{SU}(2)_L \otimes \mathrm{SU}(2)_R \otimes \mathrm{U}(1)_{B-L}$ model with $\tan\beta = 40$.

restricts $\sigma(pp \to W') \times \mathrm{BR}(W' \to WZ)$ to be less than ~ 7 fb.

One of the most popular models in which an $\mathrm{SU}(2)_L$ singlet W' appears is the so-called left-right symmetric model based on the $\mathrm{SU}(2)_L \otimes \mathrm{SU}(2)_R \otimes \mathrm{U}(1)_{B-L}$ gauge theory. In this model, the W'-quark coupling is given by the $\mathrm{SU}(2)_R$ gauge coupling constant g_R; $g_{ud} = g_R$. If right-handed neutrinos in this model are heavier than W', then W' does not decay into the right-handed neutrinos, and thus this model realizes the setup of the simplified model we have discussed here. In this model, the SM Higgs field is embedded into a bifundamental representation of $\mathrm{SU}(2)_L \otimes \mathrm{SU}(2)_R$. Once this bifundamental field acquires the VEV, the W-W' mixing is induced and given by

$$\tan 2\zeta \simeq 2\sin 2\beta \left(\frac{g_R}{g_L}\right) \frac{M_W^2}{M_{W'}^2}, \tag{7}$$

where g_L is the $\mathrm{SU}(2)_L$ gauge coupling constant and $\tan\beta$ is the ratio between the diagonal components of the bifundamental Higgs VEV. In Fig. 1, we also show the value of ζ obtained through this relation as a function of $g_{ud} = g_R$ for $\tan\beta = 40$ in the brown dashed line. This value of $\tan\beta$ is favored to explain the top and bottom quark masses in this model. It is found that although the predicted values evade the electroweak precision bound, it is far below the values required to explain the diboson excess.

Before concluding this subsection, we comment on other possible excesses reported so far which might also indicate the presence of a W' with a mass of around 2 TeV. In Ref. [21], the CMS Collaboration reported a small excess near 1.8 TeV in the searches of W' decaying into W and Higgs boson in the $l\nu bb$ final state. However, as we have seen above, this conflicts with another constraint on the $W' \to Wh$ channel given by the CMS experiment [6]. In addition, the CMS Collaboration announced a possible signal in the searches of W' decaying into the two electrons and two jets final state through a right-handed neutrino, whose peak is around 2.1 TeV with its significance being $\sim 2.8\sigma$ [22]. Though there have been several proposals for W' models that account for the 2.8σ excess [23], the models in general predict too small event rates for the diboson channel, and therefore fail to explain the ATLAS diboson resonance signal.

III. Z' MODEL

Next, we consider a leptophobic Z'. For reviews on Z' models, see Refs. [24,25]. We regard it as a gauge boson accompanied by an extra U(1) symmetry, $\mathrm{U}(1)'$, whose mass is generated after the $\mathrm{U}(1)'$ gauge symmetry is spontaneously broken. Suppose that there are two $\mathrm{SU}(2)_L$ doublets and one singlet Higgs bosons H_u, H_d, and Φ, respectively, with their $\mathrm{U}(1)'$ charges being Q'_{H_u}, Q'_{H_d}, and Q'_Φ. We further assume that H_u only couples to the up-type quarks while H_d couples to the down-type quarks and charged leptons, just as the minimal supersymmetric (SUSY) SM (MSSM) and the type-II two-Higgs-doublet model. We require $\mathrm{U}(1)'$ to be leptophobic, i.e., $Q'_L = Q'_{e^c_R} = 0$, and then this leads to $Q'_{H_d} = 0$.

After these Higgs bosons acquire VEVs, the mass matrix for the $\mathrm{U}(1)'$ gauge field $\hat{Z}'$ and a linear combination of the $\mathrm{SU}(2)_L$ and $\mathrm{U}(1)_Y$ gauge fields ($\hat{W}^a$ and $\hat{B}$, respectively), $\hat{Z} = \cos\theta_W \hat{W}^3 - \sin\theta_W \hat{B}$, is given by

$$\mathcal{L}_{\mathrm{mass}} = \frac{1}{2}(\hat{Z} \quad \hat{Z}') \begin{pmatrix} \hat{M}_Z^2 & \Delta M^2 \\ \Delta M^2 & \hat{M}_{Z'}^2 \end{pmatrix} \begin{pmatrix} \hat{Z} \\ \hat{Z}' \end{pmatrix}, \tag{8}$$

with

$$\hat{M}_Z^2 = \frac{g_Z^2 v^2}{4}, \qquad \Delta M^2 = -\frac{g_Z g_{Z'} v^2}{2} Q'_{H_u} \sin^2\beta,$$
$$\hat{M}_{Z'}^2 = g_{Z'}^2 (Q'^2_{H_u} \sin^2\beta v^2 + Q'^2_\Phi v_\Phi^2). \tag{9}$$

Here $\langle H_u^0 \rangle = v\sin\beta/\sqrt{2}$, $\langle H_d^0 \rangle = v\cos\beta/\sqrt{2}$, and $\langle \Phi \rangle = v_\Phi/\sqrt{2}$ are defined. g_Z is the gauge coupling constant given by $g_Z \equiv \sqrt{g'^2 + g^2}$ with g' and g being the $\mathrm{U}(1)_Y$ and $\mathrm{SU}(2)_L$ gauge coupling constants, respectively, and $g_{Z'}$ is the $\mathrm{U}(1)'$ gauge coupling constant. The mass eigenstates Z and Z' are then obtained through the diagonalization with an orthogonal matrix as

$$\begin{pmatrix} Z \\ Z' \end{pmatrix} = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} \hat{Z} \\ \hat{Z}' \end{pmatrix}, \qquad (10)$$

with

$$\tan 2\theta = -\frac{2\Delta M^2}{\hat{M}^2_{Z'} - \hat{M}^2_Z} \simeq 4Q'_{H_u}\sin^2\beta\left(\frac{g_{Z'}}{g_Z}\right)\frac{M_Z^2}{M_{Z'}^2}, \qquad (11)$$

where M_Z and $M_{Z'}$ are the masses of Z and Z', respectively. Again, the mixing angle is suppressed by a factor of $M_Z^2/M_{Z'}^2$.

The couplings of $\hat{Z}'$ to the SM fermions f are given by

$$\mathcal{L}_{\rm int} = g_{Z'}\bar{f}\hat{\not{Z}}'(Q'_{f_L}P_L + Q'_{f_R}P_R)f, \qquad (12)$$

with Q'_{f_L} and Q'_{f_R} being the $\mathrm{U}(1)'$ charges of the left- and right-handed components of f, respectively. Here again, we have neglected possible flavor changing effects for simplicity.

Now let us evaluate the partial decay widths of Z'. For the decay mode into quarks, we have

$$\Gamma(Z'\to q\bar{q}) = \frac{g^2_{Z'}N_C}{24\pi}M_{Z'}\left[Q'^2_{q_L} + Q'^2_{q_R} - (Q'_{q_L} - Q'_{q_R})^2\frac{m_q^2}{M^2_{Z'}}\right] \times\sqrt{1-\frac{4m_q^2}{M^2_{Z'}}}, \qquad (13)$$

where $N_C = 3$ indicates the color factor. For the $Z'\to W^+W^-$, on the other hand, we have

$$\Gamma(Z'\to W^+W^-) = \frac{g^2_{Z'}}{48\pi}Q'^2_{H_u}\sin^4\beta M_{Z'}. \qquad (14)$$

Note that this decay width again remains sizable even though the decay mode is induced via the Z-Z' mixing, since the enhancement coming from the longitudinal polarization mode compensates for the suppression factor. According to the equivalence theorem, this decay width becomes equivalent to that of the $Z'\to Zh$ mode in the decoupling limit:

$$\Gamma(Z'\to Zh) = \Gamma(Z'\to W^+W^-). \qquad (15)$$

From the above equations, we find that the decay properties of Z' are determined by the $\mathrm{U}(1)'$ charges of quarks and H_u, the $\mathrm{U}(1)'$ gauge coupling constant $g_{Z'}$, and $\tan\beta$. Among them, we can always decrease one degree of freedom via the redefinition of the $\mathrm{U}(1)'$ charge normalization. In what follows, we normalize the $\mathrm{U}(1)'$ charges such that $Q'_{H_u} = 1$. In this case, we have $Q'_{u_R} = 1 + Q'_Q$ and $Q'_{d_R} = Q'_Q$ with $Q'_Q \equiv Q'_{q_L}$, where the latter equality follows from $Q'_{H_d} = 0$.

The production cross section of Z' at the LHC8 is estimated as [24]

$$\sigma(pp\to Z') \simeq 5.2\times\left(\frac{2\Gamma(Z'\to u\bar{u}) + \Gamma(Z'\to d\bar{d})}{\rm GeV}\right)\ [\rm fb]. \qquad (16)$$

Thus, the production cross sections are determined by the quark $\mathrm{U}(1)'$ charges and $g_{Z'}$ once the Z' mass is fixed.

The production of Z' is stringently limited by the LHC experiments. For a leptophobic Z', the strong bounds come from the $Z'\to Zh$ [6], dijet [18,19], and $t\bar{t}$ [26] resonance searches. As before, we use $\sigma(pp\to Z')\times\mathrm{BR}(Z'\to Zh)\lesssim 7\,\mathrm{fb}$ and $\sigma(pp\to Z')\times\mathrm{BR}(Z'\to jj)\lesssim 100\,\mathrm{fb}$ with the acceptance $\mathcal{A}\simeq 0.6$ being assumed for the latter case. Here, $\mathrm{BR}(Z'\to jj)$ denotes the branching ratio for decaying into a pair of quark jets. The limit from the $t\bar{t}$ resonances is found to be the strongest, as we see below. The limit depends on the width of Z'. For a 2 TeV Z' with a decay width of 20 GeV, the bound is given as $\sigma(pp\to Z')\times\mathrm{BR}(Z'\to t\bar{t}) < 11$ fb, while if the decay width is 200 GeV, the bound is relaxed to be 18 fb [26].

Similarly to the case of W', the Z-Z' mixing angle is constrained by the electroweak precision measurements. For a Z' model, however, it is not appropriate to merely use the limits on the T parameter to obtain the Z-Z' bound, since the presence of Z-Z' mixing modifies the Z-boson coupling to the SM fermions simultaneously. In fact, in the case of a leptophobic Z', the constraints from the electroweak precision measurements are relaxed because Z' does not couple to electrons [27,28]. It turns out that the present limit on the mixing angle is given as $\sin\theta\lesssim 0.008$ [28], which we use in the following analysis.

In Fig. 2, we show branching ratios of the $Z'\to jj$, $Z'\to t\bar{t}$, and $Z'\to WW$ channels in the blue, green, and red lines, respectively, as functions of Q'_Q. Here, we set $M_{Z'} = 2$ TeV, and $\tan\beta = 40$ (4) for the solid (dashed) lines. The vertical gray line corresponds to the charge assignment in an E_6 inspired leptophobic Z' model often discussed in the literature [8–12], where $Q'_Q = -1/3$, $Q'_{u_R} = 2/3$, $Q'_{d_R} = -1/3$, and $Q'_\Phi = -1$. We find that the branching fraction for the diboson channel is at most ~ 0.05. The $\tan\beta$ dependence of the branching ratios is rather small; for instance, if we vary $\tan\beta$ from 40 to 4, $\mathrm{BR}(Z'\to WW)$ changes by about 10%. Then, in Fig. 3, we show a contour plot for the values of $\sigma(pp\to Z')\times\mathrm{BR}(Z'\to WW)$ as a function of $g_{Z'}$ and Q'_Q. Here, we set $M_{Z'} = 2$ TeV and $\tan\beta = 40$, and the vertical gray line shows the charge assignments in the E_6 inspired leptophobic Z' model mentioned above. The blue and dark-blue shaded regions are excluded by the resonance searches in the dijet and $Z'\to Zh$ channels, respectively. The dark (light) gray area is excluded by the $t\bar{t}$ resonance search if the Z' decay width is 200 (20) GeV. We also show the total decay width $\Gamma_{\rm tot}$ in green dashed lines. Contrary to the case of W', the electroweak precision measurements give no constraint in the parameter region shown in the figure, since

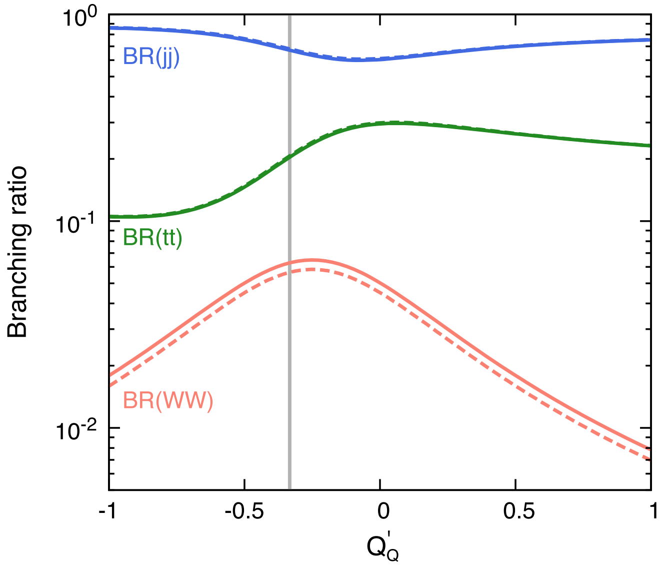

FIG. 2 (color online). Branching ratios of dijet, top quarks, and WW channels as functions of Q'_Q from top to bottom. Solid (dashed) lines represent the case of $\tan\beta = 40$ (4). We set $M_{Z'} = 2$ TeV. The vertical gray line corresponds to the charge assignment in the E_6 inspired leptophobic Z' model mentioned in the text.

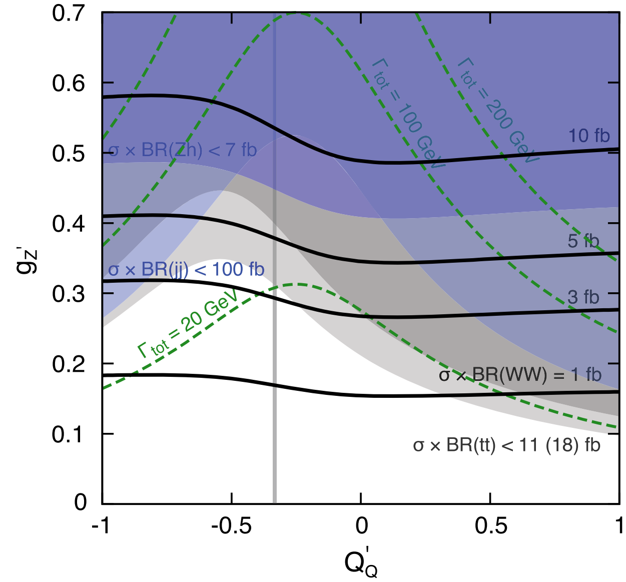

FIG. 3 (color online). Contour for the values of $\sigma(pp \to Z') \times \mathrm{BR}(Z' \to WW)$ in black solid lines. We also show the total decay width Γ_{tot} in green dashed lines. Here we set $\tan\beta = 40$ and $M_{Z'} = 2$ TeV. Blue and dark-blue shaded regions are excluded by the resonance searches in the dijet and $Z' \to Zh$ channels, respectively. The dark (light) gray area is excluded by the $t\bar{t}$ resonance search if the Z' decay width is 200 (20) GeV. The vertical gray line corresponds to the charge assignment in the E_6 inspired leptophobic Z' model mentioned in the text.

the leptophobic nature of Z' considerably weakens the limit on the Z-Z' mixing angle as discussed above. As can be seen from Fig. 3, the total decay width is well below 100 GeV in the allowed parameter region in the figure. We also find that the $t\bar{t}$ resonance search gives the most stringent constraint. In particular, if $Q'_Q = -1/3$, then $\sigma(pp \to Z') \times \mathrm{BR}(Z' \to WW)$ should be less than about 5 fb, which corresponds to $g_{Z'} \lesssim 0.4$.

Before concluding this section, we discuss an ultraviolet completion of the simplified leptophobic Z' model. In general, a leptophobic symmetry causes gauge anomaly, and hence we have to add extra $\mathrm{U}(1)'$-charged chiral fermions so that the contribution of these extra fermions removes the anomaly. A simple way to find such a set of extra fermions is to embed the SM particle content into a realization of an anomaly-free gauge group. Indeed, it turns out that GUTs based on the SUSY E_6 gauge group provide a natural framework to realize the leptophobic $\mathrm{U}(1)'$ symmetry [8–12].[2]

In SUSY E_6 GUTs, all of the MSSM matter fields as well as the right-handed neutrino superfields are embedded into a **27**-representational superfield in each generation. The **27** representation also contains new vectorlike superfields and a singlet field with respect to the SM gauge symmetry. A part of these vectorlike fields is identified as the MSSM Higgs fields. The rank of E_6 group is six, and thus this gauge group yields two additional U(1) gauge symmetries after the breaking of the GUT symmetry. It is found to be possible to construct a leptophobic U(1) charge out of a linear combination of the generators of these extra U(1) symmetries and $\mathrm{U}(1)_Y$.[3] The new states in a **27** are charged under these U(1) symmetries, which make the theory anomaly free. Among the new states, the SM singlet scalar component may acquire a VEV after the SUSY breaking, which breaks the leptophobic $\mathrm{U}(1)'$ symmetry. If the SUSY breaking occurs around the TeV scale, we expect the VEV is also $\mathcal{O}(1)$ TeV. Then the massive gauge boson associated with the spontaneous breaking could be regarded as the 2 TeV leptophobic Z' discussed in this section.

As already noted above, the $\mathrm{U}(1)'$ charge assignments for quarks are $Q'_Q = -1/3$, $Q'_{u_R} = 2/3$, and $Q'_{d_R} = -1/3$. The right-handed neutrinos and H_u have a unit charge, while the SM singlet scalar component Φ has $Q'_\Phi = -1$. By definition, the charged leptons have the zero charge, which results in $Q'_{H_d} = 0$. The masses of the extra vectorlike particles are given by the VEV of Φ. Therefore, these particles also lie around the TeV scale.

[2] For a review of the E_6 SUSY GUT models, see, e.g., Ref. [29].

[3] As discussed in Ref. [9], a kinetic mixing of the gauge fields associated with these U(1) symmetries may realize such a desirable linear combination. However, a concrete realization of the leptophobic $\mathrm{U}(1)'$ model from the E_6 GUT requires additional considerations for subtleties such as gauge coupling unification, the rapid proton-decay problem, and the structure of the Yukawa couplings. These possible issues are beyond the scope of the present paper.

This SUSY leptophobic Z' model has several phenomenologically interesting features. First, the $\mathrm{U}(1)'$ symmetry forbids the mass term for Higgsinos, i.e., the μ-term, and it is effectively induced through the Yukawa coupling between the Higgsinos and the SM singlet field which breaks the $\mathrm{U}(1)'$ symmetry, similar to the next-to-minimal SUSY SM (NMSSM). As a result, the Higgsino mass and the Z' mass have the same origin; in particular, the effective μ-parameter is expected to be around the TeV scale, which solves the so-called μ-problem. Second, in this model, there are new tree-level contributions to the Higgs mass: the F-term contribution via the singlet-Higgsino coupling just like the NMSSM and the D-term contribution of the extra $\mathrm{U}(1)'$. Taking into account these contributions as well as the one-loop correction to the charge parity (CP)-even scalars [30], we find that the observed value of the Higgs mass $\sim$125 GeV [31] is accounted for with $\mathcal{O}(1)$ TeV stops and a small value of $\tan\beta$ when $g_{Z'} = 0.4$–0.5. This should be contrasted with the MSSM prediction; in this case, if stops have masses of around 1 TeV, the observed Higgs mass is achieved with a large value of $\tan\beta$, while if $\tan\beta$ is small then stops in general have masses much larger than $\mathcal{O}(1)$ TeV to explain the Higgs mass. Notice that the 2 TeV Z' favors TeV-scale SUSY particles since the SUSY and $\mathrm{U}(1)'$ breaking scales are related to each other through the soft mass term for Φ in the scalar potential, which triggers the $\mathrm{U}(1)'$ breaking. Therefore, this model provides a natural framework for a light stop scenario without conflicting with the 125 GeV Higgs mass, which is desirable from the viewpoint of the electroweak fine-tuning problem. As various particles are predicted to have masses of $\sim$1 TeV, not only the further investigations of the diboson events, but also the direct searches of these particles in the LHC run II play an important role to test this model.[4]

IV. SUMMARY

We have considered some extensions of the SM that could explain the excesses recently reported by the ATLAS Collaboration [1]. These possible signals are found in the diboson resonance searches with two fat jets in the final state, and to account for the signals, production cross sections of $\sim$6 fb with narrow decay widths are required. These excesses may be reproduced by an extra vector boson with 2 TeV mass, and we investigated the W' and Z' models, especially. W' boson is, for example, predicted by the additional $\mathrm{SU}(2)_R$ gauge symmetry, and decays to not only the SM fermions but also W and Z bosons through the W-W' mixing. There is a tension between the excess and the bounds from the Drell-Yan processes, which forces us to forbid the leptonic decay of W' by making the right-handed neutrinos heavy. We also suffer from the tree-level flavor changing couplings of W', and thus we need to find out a way to forbid the couplings. Besides, the constraint from the electroweak precision measurements is too severe to reproduce the excess in the diboson channel. Eventually, we conclude that it is difficult to interpret the diboson signal as the W' resonance, unless the above difficulties are evaded with additional conspiracy.

Z' boson is another good candidate for the diboson resonance. It also appears in various new physics models; for instance, an extra $\mathrm{U}(1)'$ symmetry is predicted by the GUTs based on large gauge groups, in which the $\mathrm{U}(1)'$ charges are generally assigned to the SM fermions and the Higgs field according to their gauge structure. In this case, the Z' associated with this $\mathrm{U}(1)'$ symmetry decays to a pair of the SM fermions, W^+W^- and Zh through the Z-Z' mixing generated by the kinetic mixing and mass mixing due to the nonzero VEV of the $\mathrm{U}(1)'$-charged Higgs field. Again, a leptophobic Z' has an advantage to suppress the Drell-Yan bound. In Sec. III, we investigate such a possibility and find that there is a sizable parameter region that could explain a large part of the excess and is still allowed by the current experimental constraints. We also consider a concrete leptophobic Z' model inspired by the E_6 GUT, and discuss its implication on the Higgs mass and the SUSY scale. This model predicts the new vectorlike particles and the SUSY particles to have TeV-scale masses, which are accessible at the next stage of the LHC running. We therefore expect that the future LHC experiments may not only provide us a deeper understanding for the ATLAS diboson excess, but also shed light on new physics behind simplified models discussed in this paper.

Finally we briefly comment on the diboson resonance searches performed by the CMS Collaboration [32,33]. The CMS Collaboration has searched for resonances decaying into two gauge bosons in hadronic final states [32], similarly to the ATLAS search [1], as well as in semi-leptonic final states [33]. Interestingly, a small excess was found in both of these searches around 1.8 TeV for the resonance mass, which might be the same origin as the ATLAS diboson anomaly. If it is not the case, the semi-leptonic search result [33] gives a stringent limit on the 2 TeV excess observed by the ATLAS Collaboration. After all, further searches of the diboson events are indispensable for confirming or excluding the 2 TeV diboson anomaly, and are to be done in the near future.

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[4]We however note that the decay branching ratios of Z' in this model may be different from those presented above if some of the additional particles have masses smaller than 1 TeV.

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