Drell-Yan Constraints on New Electroweak States: LHC as a \( pp \rightarrow \ell^+\ell^- \) Precision Machine

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The Standard Model (SM) extensions with vector-like states which have either zero hypercharge or zero weak isospin are rather poorly constrained by the electroweak precision measurements. Such new states would however modify the running of the gauge couplings at high energies. As a result, the Drell-Yan process \( pp \rightarrow \ell^+\ell^- \) at the LHC places useful constraints on these models. The relevant observables include both the di-lepton invariant mass distribution \( M_{\ell\ell} \) and the forward-backward asymmetry \( A_{FB} \). We find that the LHC Run 1 data and the initial data from Run 2 surpass the sensitivity of LEP and already put meaningful constraints on the existence of such particles, which will become progressively stronger with more data.

INTRODUCTION

The existence of new particles at the electroweak (EW) scale can be motivated by various considerations including the hierarchy problem and the WIMP solution to the dark matter puzzle. Many such models are constrained by measurements of EW precision observables (EWPO) (see e.g. \cite{1}). However, if the new states have either zero hypercharge or zero weak isospin, such constraints become very weak. Direct searches for these states at the LHC are highly model dependent, especially if the new fermions carry no color. The current bounds can be as low as \( M_f \sim 100 - 200 \) GeV depending on the specifics of the model \cite{2,3}.

On the other hand, the presence of new EW states affects the renormalization group (RG) running of the SM gauge couplings at high energies \cite{6}, which could be probed at the LHC via precise measurements of the Drell-Yan production rate \cite{7}. In this work, we focus on the process \( pp \rightarrow \gamma^*/Z^* \rightarrow \ell^+\ell^- \) at large di-lepton invariant masses. We find that the LHC Run 1 data at \( \sqrt{s} = 8 \) TeV greatly surpass the sensitivity of EWPO and provide non-trivial constraints on the existence of such electroweak states. We also show that the LHC Run 2 data will significantly improve the LHC 8 TeV bounds. Moreover, unlike the bounds from direct searches, these limits are insensitive to the details of the spectrum and couplings among the new states.

The Drell-Yan process \( pp \rightarrow \gamma^*/Z^* \rightarrow \ell^+\ell^- \) at high di-lepton invariant mass \( M_{\ell\ell} \) combines a significant cross section \( \sigma \) with good control of theoretical and experimental uncertainties, allowing for precise measurements of the SM EW gauge couplings \cite{6}. Here we explore two observables: the di-lepton invariant mass distribution \( d\sigma/dM_{\ell\ell} \) and the di-lepton forward-backward asymmetry \( A_{FB} \). These observables are complementary, being sensitive to different combinations of the gauge couplings. We stress that \( d\sigma/dM_{\ell\ell} \) data from the 8 TeV LHC Run already place interesting constraints on the scenarios at hand, while \( A_{FB} \) will become important with more data.

DRELL-YAN SIGNATURES OF NEW PHYSICS

In order to assess the sensitivity of \( d\sigma/dM_{\ell\ell} \) and \( A_{FB} \) to new physics, let us extend the SM by a set of vector-like fermions with mass \( M_f \). To isolate their effect on \( \alpha_Y = g_Y^2/4\pi \) and \( \alpha_2 = g_2^2/4\pi \), respectively, we assume that the new fermions are either \( N \) copies of \( SU(2)_L \) singlets with hypercharge \( Y_f \), or \( N_q \) copies of \( SU(2)_L \) doublets with zero hypercharge (here \( N, N_q \) include the color factor for the case of vector-like quarks). Such SM extensions are very weakly constrained by EWPO. The relevant EW parameters are \( Y = N Y_f^2 \times \alpha_Y m^2_W/(15\pi M_f^2) \) and \( W = N_q \times \alpha_2 m^2_W/(30\pi M_f^2) \), which are constrained at the level of a few per mille \cite{1}, rendering the EWPO constraints essentially irrelevant for \( M_f > 400 \) GeV.

The new states contribute to the RG running of the EW couplings by changing the SM beta functions by

\[
\Delta \beta_Y = \frac{g_Y^3}{16\pi^2} \frac{4N Y_f^2}{3}, \quad \Delta \beta_2 = \frac{g_2^3}{16\pi^2} \frac{2N_q}{3}
\]

for scales \( \mu > M_f \). In what follows, we focus on the leading log corrections to \( \alpha_i \) at high energies \( \mu \gg M_Z \), \( \delta\alpha_i(\mu) \equiv \alpha_i(\mu) - \alpha_i(M_Y) \).

Di-Lepton Invariant Mass Distribution. The Born level \( pp \rightarrow \gamma^*/Z^* \rightarrow \ell^+\ell^- \) differential cross section is given by

\[
\frac{d\sigma}{dM_{\ell\ell}} = \frac{\pi}{648 M_{\ell\ell}^3} \sum_q C_q(\alpha_Y, \alpha_2, M_{\ell\ell})
\]

\[
\times \int_{\ln\left(\frac{M_{\ell\ell}}{M_Z}\right)}^{\ln\left(\frac{\sqrt{s}}{M_{\ell\ell}}\right)} dy \left[x_1 f_1(x_1) x_2 f_2(x_2)\right],
\]

where \( M_{\ell\ell} \) is the di-lepton invariant mass, \( \sqrt{s} \) is the center-of-mass energy of the pp collision, \( f_i \) are the parton distribution functions (PDFs), \( x_{1,2} = \frac{M_{\ell\ell}}{\sqrt{s}} e^{\pm y} \) are the Bjorken variables and \( y \) is the di-lepton rapidity (see \cite{8,10}). The \( C_u \) and \( C_g \) coefficients are given by (neglecting the Z boson width \( \Gamma_Z \))
\[ C_u(\alpha_Y, \alpha_2, M_{\ell\ell}) = 576 \alpha_Y^2 \alpha_2^2 \times \left[ \frac{4}{9} - \frac{\frac{1}{2} - \frac{4}{2} s_W^2}{3 s_W^2 (1 - s_W^2)} \left( \frac{1}{2} + \frac{2}{2} s_W^2 \right) \right] + \frac{\frac{4}{9} + \frac{\frac{1}{2} - \frac{2}{2} s_W^2}{6 s_W^2 (1 - s_W^2)}}{16 s_W^2 (1 - s_W^2)^2} \] \\
\[ C_d(\alpha_Y, \alpha_2, M_{\ell\ell}) = 576 \alpha_Y^2 \alpha_2^2 \times \left[ \frac{1}{9} + \frac{\frac{1}{2} - \frac{2}{2} s_W^2}{6 s_W^2 (1 - s_W^2)} \left( \frac{1}{2} - \frac{2}{2} s_W^2 \right) \right] + \frac{\frac{1}{9} + \frac{\frac{1}{2} - \frac{2}{2} s_W^2}{16 s_W^2 (1 - s_W^2)^2}}{16 s_W^2 (1 - s_W^2)^2} \] 

with \( s_W^2 = \alpha_Y/(\alpha_2 + \alpha_Y) \). In the limit \( M_{\ell\ell} \gg m_Z \), (3) simplifies to

\[ C_u(\alpha_1) = 85 \alpha_Y^2 + 6 \alpha_Y \alpha_2 + 9 \alpha_2^2, \]
\[ C_d(\alpha_1) = 25 \alpha_Y^2 - 6 \alpha_Y \alpha_2 + 9 \alpha_2^2. \] 

We note that this particularly simple form can be viewed as a result of the exchange of the hypercharge and \( W_3 \) bosons, which are effectively massless in this limit.

In the leading log approximation, the gauge couplings are to be understood as the running couplings taken at the scale \( M_{\ell\ell} \), which controls the partonic center-of-mass energy of the Drell-Yan process, \( \alpha_i \rightarrow \alpha_i(M_{\ell\ell}) \). In this case, (2) and (3) allow us to derive the variation of \( d\sigma/dM_{\ell\ell} \) under the change of the SM gauge couplings at the scale \( \mu = M_{\ell\ell} \). For instance, taking \( 1/3 = 8 \) TeV and \( \mu = 1 \) TeV, we find

\[ \frac{\delta d\sigma/dM_{\ell\ell}}{dM_{\ell\ell}} \bigg|_{\mu=1 \text{ TeV}} \approx 0.84 \frac{\delta \alpha_Y(\mu)}{\alpha_Y(\mu)} + 1.16 \frac{\delta \alpha_2(\mu)}{\alpha_2(\mu)}, \] 

using NNPDF2.3 [11] at Next-to-Next-to-Leading-Order (NNLO). This illustrates that an \( (1-10\%) \) precision measurement of \( d\sigma/dM_{\ell\ell} \) can be translated into a bound on the variation of \( \delta \alpha_Y/\alpha_Y(\mu) \) or \( \delta \alpha_2/\alpha_2(\mu) \) of the same order. In order to assess the sensitivity to new states, we define the ratio \( R(\mu, M_f, \Upsilon) = \frac{\sigma(\mu, M_f, \Upsilon)}{dM_{\ell\ell}}/dM_{\ell\ell}(\mu) \) with \( \Upsilon = NY_f^2 \) (\( \Upsilon = N_d \)) for the (SU(2)\_L) singlet (doublet) vector-like fermions.

We now turn to deriving constraints on the new EW states using the NNLO Drell-Yan cross section [12]. The data include e\(^+\)e\(^-\) and \( \mu^+\mu^- \) lepton pairs with \( M_{\ell\ell} \) up to 2 TeV, and the agreement between the SM prediction and experimental data is excellent, with the error bars well below 10\% all the way up to 1 TeV. We use the CMS data in the binned interval \( M_{\ell\ell} \in [220 \text{ GeV}, 2 \text{ TeV}] \). Since the bins have a sizable \( \Delta M_{\ell\ell} \), the ratio \( R(\mu_i, M_f, \Upsilon) \) should be defined for each bin \( i \) as an integral of \( d\sigma/dM_{\ell\ell} \) over \( \mu \) in the range \( \Delta M_{\ell\ell} \), divided by the corresponding integral of \( d\sigma_{\text{SM}}/dM_{\ell\ell} \) over the same \( \mu \) range. We note that although (2) is defined at LO in the hard scattering process, we include corrections up to NNLO through a \( \mu \)-dependent rescaling of the LO cross section. To this end, we use the NNLO theoretical prediction for the Drell-Yan cross section [12] via FEWZ 3.1 [13] [14]. We then perform a \( \chi^2 \) fit to the binned CMS 8 TeV data using NNPDF2.3 [11] NNLO PDFs (and retaining finite \( m_Z \) effects) with

\[ \chi^2(M_f, \Upsilon) = \sum_{\mu_i} \left( \frac{R_{\text{exp}}(\mu_i, M_f, \Upsilon) - R(\mu_i, M_f, \Upsilon)}{\Delta R_{\text{exp}}(\mu_i, M_f, \Upsilon)} \right)^2. \] 

The data to theory ratio \( R_{\text{exp}} \) for each bin as well as the corresponding uncertainty \( \Delta R_{\text{exp}} \) are extracted from the CMS 8 TeV analysis [12]. We show the resulting 1\( \sigma \) and 2\( \sigma \) bounds on \( NY_f^2 \) vs \( M_f \) and \( N_d \) vs \( M_f \) in Figure 1 Left and -Right, respectively. The constraints on \( NY_f^2 \) and \( N_d \) are seen to be much stronger than those from EWPO. In Figure 1 Left, we also show the bound from the SM gauge coupling perturbativity up to \( \mu = 3 \) TeV.

We go on to compute the projected LHC limits for \( \sqrt{s} = 14 \) TeV with 300 fb\(^{-1}\), under the assumption that future experimental data agree with the SM prediction (\( R_{\text{exp}} = 1 \)). The largely dominant source of systematic uncertainties in the Drell-Yan process for \( M_{\ell\ell} \gg m_Z \) at \( \sqrt{s} = 14 \) TeV is given by the PDF uncertainty [6] (see also [15]). We use the accurate estimate of the size of the PDF uncertainties in the Drell-Yan process from [6], adding this uncertainty in quadrature with the statistical uncertainty (which becomes dominant for \( \mu \gtrsim 1 \) TeV) to obtain a projected estimate of \( \Delta R_{\text{exp}} \). We then perform a \( \chi^2 \) fit using (6), and choosing \( \mu_i = 1, 1.5, 2, 2.5, 3 \) TeV. Our motivation for this choice is twofold: A substantial \( M_{\ell\ell} \) separation between the \( \mu_i \) bins implies weak correlations among them, allowing us to treat the bins in [6] as uncorrelated.

In addition, for values of \( \mu \) beyond 3 TeV, the statistical uncertainty in the Drell-Yan process with 300 fb\(^{-1}\) becomes \( O(1) \). Our 1\( \sigma \) and 2\( \sigma \) projected bounds on \( NY_f^2 \) vs \( M_f \) (\( N_d \) vs \( M_f \)) with \( \sqrt{s} = 14 \) TeV and 300 fb\(^{-1}\) are shown in Figure 1 Left (Figure 1 Right).

Before moving on, a few comments are in order: (i) The models considered here are idealized in that all the new fields are taken to be identical. In reality one may expect a more complicated mass spectrum, and/or several states with different EW quantum numbers. In this respect, the presence of states with other quantum numbers could affect dramatically the collider phenomenology of these models leading for example to complicated

\footnote{Our results agree quantitatively with those in [6] [15], where these weak correlations have been included.}
final states which would obscure LHC searches. On the other hand, these states do not necessarily affect the SM gauge coupling running in any tangible way. This highlights the fact that our bounds rely on the cumulative effect of new EW states and are fairly insensitive to further details of the model. (ii) Employing the leading log approximation for the SM gauge couplings is a source of theory uncertainty. However, in our 8 TeV analysis this error is far subleading to the PDF, background and statistical uncertainties, and can safely be neglected. In the 14 TeV analysis, this source of error is negligible for the $M_f$ mass range considered here, except for $M_f \rightarrow 1$ TeV where it is again largely subdominant to the PDF and statistical uncertainties. (iii) Besides the neutral current Drell-Yan process $pp \rightarrow \gamma^*/Z^* \rightarrow \ell^+\ell^-$, the charged current process $pp \rightarrow W^* \rightarrow \ell\nu$ can also be used to constrain $\alpha_2$ via the transverse mass distribution $m_T [6]$. This analysis is beyond the scope of our current study. In the next Section, we show that the forward-backward asymmetry can yield useful information complementary to that from the invariant mass distribution.

\[ \frac{d\sigma(q\bar{q} \rightarrow \ell^+\ell^-)}{d\cos \theta^*} = \frac{1}{3} \left[ (1 + \cos^2 \theta^*) F_0^q(M_{\ell\ell}^2) + 2 \cos \theta^* F_1^q(M_{\ell\ell}^2) \right], \quad (7) \]

where $\theta^*$ is the angle between the quark and $\ell^-$ momenta.

Di-Lepton Forward-Backward Asymmetry. In addition to the di-lepton invariant mass distribution $d\sigma/dM_{\ell\ell}$, an observable which can be useful in constraining the gauge couplings is the forward-backward asymmetry $A_{\text{FB}}$. At the parton level, in the di-lepton center-of-mass frame, one has [16]

\[ A_{\text{FB}} = \frac{1}{2} \left[ \frac{d\sigma(q\bar{q} \rightarrow \ell^+\ell^-)}{d\cos \theta^*} \right]_{\text{forward}} - \left[ \frac{d\sigma(q\bar{q} \rightarrow \ell^+\ell^-)}{d\cos \theta^*} \right]_{\text{backward}} \]

where $\theta^*$ is the angle between the quark and $\ell^-$ momenta.

For $M_{\ell\ell} \gg m_Z$, the form–factors $F_{0,1}^q(M_{\ell\ell}^2)$ read

\[ F_0^q = \frac{\pi C_4(\alpha_s)}{1152 M_{\ell\ell}^2}, \quad F_1^q = \frac{\pi (15\alpha_s^2 + 2\alpha_s \alpha_2 + 3\alpha_2^2)}{384 M_{\ell\ell}^2}, \]

\[ F_0^d = \frac{\pi C_4(\alpha_s)}{1152 M_{\ell\ell}^2}, \quad F_1^d = \frac{\pi (3\alpha_s^2 - 2\alpha_s \alpha_2 + 3\alpha_2^2)}{384 M_{\ell\ell}^2}. \]

The forward-backward asymmetry $A_{\text{FB}}$ is defined by

\[ A_{\text{FB}}(M_{\ell\ell}^2) = \frac{\sigma(\cos \theta^* > 0) - \sigma(\cos \theta^* < 0)}{\sigma(\cos \theta^* > 0) + \sigma(\cos \theta^* < 0)} = \frac{3F_0^q(M_{\ell\ell}^2)}{4F_0^q(M_{\ell\ell}^2)}. \]

It is proportional to the quark and lepton axial couplings. In order to understand the sensitivity of $A_{\text{FB}}$ to the gauge

![Graph showing the dependence of the forward-backward asymmetry on the invariant mass $M_{\ell\ell}$](image-url)

**FIG. 1.** Left: 8 TeV $(19.7 \text{ fb}^{-1})$ Drell-Yan $1\sigma$ (orange) and $2\sigma$ (red) bounds on $SU(2)_L$ singlet vector-like fermions with hypercharge $Y_f$ and multiplicity $N$ vs mass $M_f$. The dashed red (orange) lines show the projected $2\sigma$ ($1\sigma$) bounds from 14 TeV LHC $(300 \text{ fb}^{-1})$. Also displayed are the EWPO constraint from the $Y$-parameter $Y < 2 \times 10^{-3}$ (brown) and the $g_\pi$ perturbativity constraint up to the scale $\mu \sim 3 \text{ TeV}$, $g_\pi^2 (3 \text{ TeV}) < 4\pi$ (light-blue). Right: Same for $SU(2)_L$ doublets with zero hypercharge and multiplicity $N_d$. The EWPO constraint is due to the W-parameter $W < 1 \times 10^{-3}$. 


couplings, let us consider the parton level variations

\[
\frac{\delta A^u_{FB}}{A^u_{FB}} \bigg|_{\mu=1 \text{ TeV}} \simeq 0.27 \left( \frac{-\delta \alpha_Y(\mu)}{\alpha^S_Y(\mu)} + \frac{\delta \alpha_L(\mu)}{\alpha^S_L(\mu)} \right),
\]

\[
\frac{\delta A^d_{FB}}{A^d_{FB}} \bigg|_{\mu=1 \text{ TeV}} \simeq 0.35 \left( \frac{-\delta \alpha_Y(\mu)}{\alpha^S_Y(\mu)} + \frac{\delta \alpha_L(\mu)}{\alpha^S_L(\mu)} \right),
\]

where in the SM \(A^u_{FB} \simeq 0.6\) at \(\mu = 1\) TeV. Comparing this to \([3]\), we immediately see that \(A_{FB}\) can provide information complementary to that from \(d\sigma/dM_{\ell\ell}\). Although the asymmetry is less sensitive to the variation of the couplings than the differential cross section, it can still be a useful observable, having the advantage that many systematic and QCD uncertainties cancel in \(A_{FB}\). It has however its own problematic issues, above all the reconstruction of the quark direction, which is reliable for large rapidities only \([16]\). Current measurements at 8 TeV LHC \([17, 18]\) contain substantial uncertainties of \(\sim 700\) GeV for \(|Q| = 2e\) to \(|Q| = 8e\) \([20]\). However, if they decay promptly to multi-lepton final states and/or jets, no general bounds exist as these are highly model dependent and would require a dedicated analysis.

For VL leptons with exotic, non-SM charges we can only decay via higher-dimensional operators. If these are stable on detector scales, the lower limit on their masses is \(M_f \gtrsim 700\) . . . 800 GeV for \(|Q| = 2e\) to \(|Q| = 8e\) \([20]\). However, if they decay promptly to multi-lepton final states and/or jets, no general bounds exist as these are highly model dependent and would require a dedicated analysis.

CONCLUSIONS

Current and future Drell-Yan measurements at the LHC place important constraints on possible new states with EW quantum numbers. Such states modify the RG running of the SM gauge couplings at high energies thereby affecting the Drell-Yan production rates. The resulting bounds are fairly insensitive to the details of the mass spectrum or interactions among these states and therefore are complementary to the direct search constraints. We have shown that both the di-lepton differential \(d\sigma/dM_{\ell\ell}\) distribution and the forward-backward asymmetry \(A_{FB}\) can place meaningful constraints on the existence of new EW states, and that the LHC sensitivity of \(pp \rightarrow \ell^+\ell^-\) to the new states far surpasses that of the current EW precision measurements.

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