Tevatron $Wjj$ Anomaly for a Model with Two Different Mechanisms for Mass Generation of Gauge Fields

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Abstract The latest Fermilab Collider Detector (CDF) anomaly, namely the excess of dijet events in the invariant-mass window 120–160 GeV in associated production with a $W$ boson, is explained by a baryonic new neutral vector $C$-boson, of mass (145 GeV), predicted by the Wu mechanisms for mass generation of gauge field. The Standard Model (SM) $W$, $Z$-bosons normally get their masses through the coupling with the SM Higgs particle of mass 114–200 GeV. Here, the baryonic $C$-boson has negligible couplings with leptons and, thus, is unaffected by the dilepton $C$ constraints.

Keywords Collider Detector Fermilab · Electroweak · Interactions · Gauge field · Symmetry breaking

1 Introduction

The latest surprise brought forth by the Fermilab’s Collider Detector (CDF) experiment is the $Wjj$ production: an excess of the number of events in the invariant-mass window of 120–160 GeV in the dijet system of the associated production of a $W$-boson with 2 jets [1–4].

The hypothesis behind Fermilab’s CDF experiment predicted that the number of the events—producing a $W$ boson and a pair of jets—would fall off as the mass of the jet pair increased [1–4]. The CDF experimental data, however, showed something strange (see [2]): a bump in the number of events when the mass of the jet pair was about 145 GeV. The excess of events in the window $M_{jj} \sim 120–160$ GeV appears to be a resonance. It is reportedly at 4.8 sigma, tantalizingly close to five-sigma certainty [2, 3].

However, DZero (D0) Collaboration [27], have cross-checked the observation with their own independent data and analysis tools, and have found no evidence of a new particle. Their detectors are somewhat different in design and in the selection criteria used to analyze events. It is certainly possible that CDF has seen something that DZero has missed.
From the distribution of events in the CDF experiment we can see that the width of the resonance appears to be slightly wider than the SM \( Z \)-boson.

The candidate particle may not belong to the standard model of particle physics. Instead, as some argue [5–7], it might be the first hint of a new force of nature, termed technicolour. This force would resolve some problems with the Standard Model (SM), such as the “naturalness problem” [28]; however, it predicts the existence of many light particles that should have been detected by now [28].

Most workers agree that the mysterious particle produced by the CDF experiment is not the long-sought Higgs boson [8–15], believed by many to endow particles with mass. If it were, the bump in the experimental data would be 300 times smaller. What’s more, a Higgs particle would probably decay into bottom quarks, which seem to be absent from the Fermilab data [1–4].

In this paper, we investigate the nature of this mysterious particle and propose a model with two different mechanisms for mass generation of gauge fields. The Tevatron \( W_{ij} \) anomaly is explained by a baryonic new neutral vector \( C \)-boson of mass (145 GeV), as predicted by the Wu mechanisms for mass generation of gauge fields [16–19, 30, 31]. The SM \( W, Z \)-bosons normally acquire their masses through their coupling with the SM Higgs boson, mass 114–200 GeV [20–23]. The baryonic \( C \) boson has negligible couplings to leptons, and so is not affected by the dilepton \( C \) constraints.

2 The Lagrangian of the Model

Suppose that the gauge symmetry of the theory is \( SU(N) \times U(2) \) group, which is written specifically as follows:

\[
G = SU(N) \times U(2),
\]

where \( SU(N) \) is the special unitary group of \( N \)-dimensions, \( \psi(x) \) is a \( N \)-component vector in the fundamental representative space of \( SU(N) \) group, and \( T_i \) \((i = 1, 2, \ldots, N^2 - 1)\) denotes the representative matrices of the generators of \( SU(N) \) group. The latter are Hermit and traceless. They satisfy the condition:

\[
[T_i, T_j] = \text{if}_{ijk} T_k, \quad \text{Tr}(T_i T_j) = \delta_{ij} K
\]

where \( \text{if}_{ijk} \) are structure constants of the \( SU(N) \) group, and \( K \) is a constant independent of the indices \( i \) and \( j \) but dependent on the representation of the group. The representative matrix of a general element of the \( SU(N) \) group is expressed as:

\[
U = e^{i\alpha^i T_i},
\]

with \( \alpha^i \) being the real group parameters. In global gauge transformations, all \( \alpha^i \) are independent of space-time coordinates, while in local gauge transformations \( \alpha^i \) are functions of space-time coordinates. \( U \) is a unitary \( N \times N \) matrix.

In order to introduce the mass term of gauge fields without violating local gauge symmetry at energy scale close to 2 TeV, two kinds of gauge fields are required: \( a_\mu \) and \( b_\mu \) [16].

In this version of Wu gauge model, the gauge fields \( a_\mu^i \) and \( b_\mu^i \) are introduced as follows. Let the Lagrangian of our theory to be invariant under a set of infinitesimal transformations of the spinor fields \( \psi(x) \) [32].

\[
\delta \psi(x) = i \alpha^i(x) T_i \psi(x),
\]
with $T_i$ being some of the independent constant matrices, and $\alpha^i(x)$ the real infinitesimal parameters.

Differentiating Eq. (4) we derive

$$
\delta\left(\partial_\mu \psi(x)\right) = i\alpha^i(x) T_i \left(\partial_\mu \psi(x)\right) + i \left(\partial_\mu \alpha^i(x)\right) T_i \psi(x).
$$

(5)

To make the Lagrangian invariant, we artificially introduce the field $a^i_\mu$ whose transformation rule involves a term $\partial_\mu \alpha^i$. The latter term can be used to cancel the second term in Eq. (5).

Here, the other kind of gauge fields, $b^i_\mu$, is generated due to nonsmoothness of the field trajectories for the scalar phases of the spinor fields in the $SU(N)$ gauge theory. This is achieved in the following steps:

First, let us consider a Lagrangian for free spinor fields:

$$
L = i \bar{\psi}^j \gamma^\mu \partial_\mu \psi^j - m \bar{\psi}^j \psi^j,
$$

(6)

where $j = 1, 2, \ldots, N$ (in what follows the index $j$ will be omitted).

The Lagrangian (6) is invariant under global non-Abelian $SU(N)$ transformations

$$
\delta \psi(x) = e^{iT_i a^i} \psi(x), \quad \delta \bar{\psi}(x) = e^{-iT_i a^i} \bar{\psi}(x).
$$

(7)

In the Feynman formulation of quantum field theory, the transition amplitudes are expressed by the path integrals that are centered on nonsmooth field trajectories [33]:

$$
\langle \Phi_1, t_1 | \Phi_2, t_2 \rangle = N \int_{\Phi_1}^{\Phi_2} (D\Phi) \exp \left[ \frac{i}{\hbar} \int_{t_1}^{t_2} d^4 x L(\Phi, \partial \Phi) \right].
$$

(8)

In this context, the Lagrangian (6) and its symmetries are determined on the class of nonsmooth functions $\psi(x)$, corresponding to nonsmooth trajectories in path integrals. In the strict sense, the derivatives involved in the Lagrangian (6) are discontinuous functions. From the standpoint of physics, field trajectory nonsmoothness is related to fluctuations of the local fields. As a rule, Feynman integrals are additionally specified by the implicit switch to “smoothed-out” approximations [34]. In this case the degrees of freedom corresponding to gauge vector fields are lost. Here we show that, as in quantum electrodynamics [35], in the non-Abelian $SU(N)$ gauge theory these degrees of freedom can be explicitly taken into account when “smoothening” of nonsmooth fields is carried out carefully [36]. Let us approximate nonsmooth functions $\theta^i(x)$ by smooth functions $\alpha^i(x)$:

$$
\theta^i(x) = \alpha^i(x) + \cdots.
$$

(9)

In order to write down the next term of the “smoothed-out” representation of the nonsmooth functions $\theta^i(x)$, it is necessary to consider the behaviour of the first derivatives of $\theta^i(x)$. The derivatives $\partial_\mu \theta^i(x)$ at nonsmoothness points of $\theta^i(x)$ are discontinuous functions. Since the derivatives $\partial_\mu \alpha^i(x)$ are continuous functions, they are a poor approximation of the behaviour of the derivatives of the “smoothed-out” $\theta^i(x)$. We denote the difference between them by $\theta^{i\mu}(x)$ and write $\partial_\mu \theta^i(x)$ as follows:

$$
\partial_\mu \theta^i(x) = \partial_\mu \alpha^i(x) + \theta^{i\mu}(x).
$$

(10)

Since the nonsmooth fields $\theta^{i\mu}(x)$ do not reduce to gradients of smooth scalar fields, they are the nontrival vector fields that give rise to nonzero field strengths:

$$
\partial_\mu \theta^{i\mu}(x) = \partial_\mu \alpha^{i\mu}(x) \neq 0.
$$

(11)

Therefore, the fields $\partial_\mu \theta^i(x)$ involve those additional degrees of freedom which are related to the nonsmoothness of the $\theta^i(x)$. It should be noted that the determination of fields $\theta^{i\mu}(x)$ is ambiguous, due to the arbitrariness of choice of $\alpha^i(x)$. 

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Let us now consider $\theta_i(x)$ as scalar phases of the spinor fields $\psi(x)$, realizing the fundamental representation of the $SU(N)$ gauge group. An explicit expression of these phase degrees of freedom is:

$$\psi(x) = e^{i T^i \theta_i(x)} \psi_0(x),$$

(12)

where the spinor fields $\psi_0$ are representatives of the class of gauge equivalent fields [37], and $e^{i T^i \theta_i(x)}$ is a unitary $N \times N$ matrix. Then, provided that the Lagrangian (6) is determined on the class of nonsmooth functions $\psi(x)$, using Eq. (12) we obtain:

$$L = i \bar{\psi}_0 \gamma^\mu \partial_\mu \psi_0 + i \bar{\psi}_0 e^{-iT^i \theta_i(x)} \gamma^\mu (\partial_\mu e^{iT^i \theta_i(x)}) \psi_0 - m \bar{\psi}_0 \psi_0.$$

(13)

The matrix $e^{iT^i \theta_i}$ can be represented as a superposition of unit matrix $I$ and $SU(N)$ group generators

$$e^{iT^i \theta_i(x)} = CI + iS_i T^i.$$

(14)

Since $T_i$ are traceless matrices normalized by $\text{Tr}(T_i T_j) = \frac{1}{2} \delta_{ij}$, the coefficients $C$ and $S_i$ in Eq. (9) are given by:

$$C = \frac{1}{N} \text{Tr}(e^{iT^i \theta_i}), \quad S_i = -2i \text{Tr}(T^i e^{iT^j \theta_j}).$$

(15)

It is easy to verify that $\text{Tr}(e^{-iT^i \theta_i} \partial_\mu e^{iT^j \theta_j}) = 0$.

Then, taking into account the commutation rules for $SU(N)$ group generators [38] we can write down:

$$\left( e^{-iT^i \theta_i} \partial_\mu e^{iT^j \theta_j} \right) = i T^j b^j_\mu,$$

(16)

$$b^j_\mu = \bar{C} \partial_\mu S^j - \bar{S}^j \partial_\mu C + (f^{ijk} - id^{ijk}) \bar{S}_j S_k,$$

(17)

where $d^{ijk}$ ($f^{ijk}$) are totally symmetric (antisymmetric) structural constants of $SU(N)$ group (the over line denotes complex conjugation). Since the matrix $e^{iT^i \theta_i}$ is unitary, the following equation is valid:

$$\bar{C} S_i - \bar{S} C + (f^{ijk} - i d^{ijk}) \bar{S}_j S_k = 0.$$

(18)

Differentiating the left and right sides of Eq. (18) and using the property of antisymmetry of $f_{ijk}$, we derive:

$$b^i_\mu - \bar{b}^i_\mu = 0.$$

(19)

From this it follows that the expression (17) is a real function. Thus $b^i_\mu$ can be identified with the gauge fields. Unlike the gauge field in electrodynamics, these fields are nonlinear functions of $\theta^i(x)$. As a consequence of the nonsmoothness of the phases $\theta^i(x)$ the fields $b^i_\mu$ are also nonsmooth. If we take into account only the first term in the right hand side of relation (10), we find that the fields $b^i_\mu$ do not contribute to the dynamics, as in classical field theory [39], and that the degrees of freedom corresponding to the gauge vector fields are lost. This account of $\theta^i_\mu(x)$ enables us to interpret the fields $b^i_\mu$ as nontrivial vector fields that give rise to nonzero field strengths:

$$\partial_\mu b^i_\nu(x) - \partial_\nu b^i_\mu(x) \neq 0.$$

(20)

Let us now obtain the transformation law for the vector fields of Eq. (12). For this purpose we consider the infinitesimally smooth local transformations for the spinor fields:

$$\delta \psi_0(x) = e^{iT^i_\alpha} \psi_0(x), \quad \delta \bar{\psi}_0(x) = e^{-iT^i_\alpha} \bar{\psi}_0(x).$$

(21)
Then the Lagrangian (13) can be written as:
\[ L = i \bar{\psi}_0 \gamma^\mu \partial_\mu \psi_0 + i \bar{\psi}_0 e^{-iT^I \theta_I(x)} \gamma^\mu \left( \partial_\mu e^{iT^I \theta_I(x)} \right) \psi_0 - m \bar{\psi}_0 \psi_0. \] (22)

By defining the gauge fields \( b^i_\mu \) similarly to Eqs. (16) and (17) as:
\[ iT_i b^i_\mu(x) = e^{iT^I a_I} e^{-iT^I \theta_I(x)} \partial_\mu \left( e^{iT^T k \theta_k(x)} e^{-iT^T a_I} \right), \] (23)
we find that the transformed gauge fields \( b^i_\mu \) are related to the fields (17) as follows:
\[ b^i_\mu(x) = b^i_\mu(x) - \partial_\mu \alpha^i(x) - f_{ijk} \alpha^j(x) b^k_\mu. \] (24)

Hence, in the framework of the considered scheme of gauge field generation we derive the usual transformation law for the \( SU(N) \) gauge fields, without the local gauge invariance of the Lagrangian (13) being necessary [36].

Let us now discuss the results thus obtained. The gauge field is \( a^i_\mu \), introduced for the purpose of ensuring the local gauge invariance of the theory. The generation of the gauge field \( b^i_\mu \) is a purely quantum phenomenon. The vector gauge field \( b^i_\mu \) is generated through nonsmoothness of the scalar phase of the fundamental spinor fields.

From the viewpoint of the described scheme of the gauge field \( b^i_\mu \) generation, the gauge principle is an “automatic” consequence of the field trajectory nonsmoothness in the Feynman path integral.

\( a_\mu(x) \) and \( b_\mu(x) \) are vectors in the canonical representative space of \( SU(N) \) group. They can be expressed as linear combinations of generators, as follows:
\[ a_\mu(x) = a^i_\mu(x) T_i, \] (25)
\[ b_\mu(x) = b^i_\mu(x) T_i, \] (26)
where \( a^i_\mu(x) \) and \( b^i_\mu(x) \) are component fields of the gauge fields \( a_\mu(x) \) and \( b_\mu(x) \), respectively.

Corresponding to these two kinds of gauge fields, there are two kinds of gauge covariant derivatives:
\[ D_\mu = \partial_\mu - ig a_\mu, \] (27)
\[ D_{b_\mu} = \partial_\mu + ig b_\mu. \] (28)

The strengths of gauge fields \( a_\mu(x) \) and \( b_\mu(x) \) are defined as
\[ a_{\mu\nu} = \frac{1}{-ig} [D_\mu, D_\nu] = \partial_\mu a_\nu - \partial_\nu a_\mu - ig [a_\mu, a_\nu], \] (29)
\[ b_{\mu\nu} = \frac{1}{ig} [D_{b_\mu}, D_{b_\nu}] = \partial_\mu b_\nu - \partial_\nu b_\mu + ig [b_\mu, b_\nu] \] (30)
respectively.

Similarly, \( a_\mu(x) \) and \( b_\mu(x) \) can also be expressed as linear combinations of generators:
\[ a_{\mu\nu} = a^i_{\mu\nu} T_i, \] (31)
\[ b_{\mu\nu} = b^i_{\mu\nu} T_i. \] (32)

Using relations (2) and (29), (30), we obtain
\[ a^i_{\mu\nu} = \partial_\mu a^i_\nu - \partial_\nu a^i_\mu + g_2 f^{ijk} a^j_\mu a^k_\nu, \] (33)
\[ b^i_{\mu\nu} = \partial_\mu b^i_\nu - \partial_\nu b^i_\mu - cg_2 f^{ijk} b^j_\mu b^k_\nu. \] (34)
The Lagrangian density of the model is
\[ \mathcal{L}_{Wu} = -\bar{\psi} (\gamma^\mu D_\mu + m) \psi - \frac{1}{4K} \text{Tr}(a^{\mu\nu} a_{\mu\nu}) - \frac{1}{4K} \text{Tr}(b^{\mu\nu} b_{\mu\nu}) - \frac{\mu^2}{2K(1+c^2)} \text{Tr}\left[(a_\mu + cb_\mu)(a_\mu + cb_\mu)\right] \]  
(35)

where \( c \) is a constant.

The space-time metric is selected as \( \eta_{\mu\nu} = \text{diag}(-1, 1, 1, 1) \), \((\mu, \nu = 0, 1, 2, 3)\). According to relation (2), the above Lagrangian density \( L \) can be rewritten as:
\[ \mathcal{L}_{Wu} = -\bar{\psi} \left( \gamma^\mu (\partial_\mu - ig_2 a_\mu T_i) + m \right) \psi - \frac{1}{4} a^{\mu\nu} a_{\mu\nu} - \frac{1}{4} b^{\mu\nu} b_{\mu\nu} - \frac{\mu^2}{2(1+c^2)} (a_\mu + cb_\mu)(a_\mu + cb_\mu) \]  
(36)

This Lagrangian has strict local gauge symmetry \([16]\).

In (1), the gauge group \( SU(2)_L \times U(1)_Y \) is the known SM of electroweak (EW) interactions \([20–23]\). The generators of \( SU(2)_L \) correspond to the three components of weak isospin \( T_a \) \((a = 1, 2, 3)\). The \( U(1)_Y \) generator corresponds to the weak hypercharge \( Y \). These are related to the electric charge by \( Q = \frac{T_3}{2} + Y \).

The \( SU(2)_L \times U(1)_Y \) invariant Lagrangian is given as follows:
\[ \mathcal{L}_{EW} = -\bar{\psi} \gamma^\mu \left[ \partial_\mu - i\frac{g_1}{2} \tau \cdot A_\mu + ig_1^{'} \frac{1}{2} B_\mu \right] \psi + \bar{e}_R \gamma^\mu \left[ \partial_\mu + ig_1^{'} B_\mu \right] e_R - \frac{1}{4} F_{\mu\nu}^\alpha F^{\alpha\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu}, \]  
(37)

with the field strength tensors
\[ F_{\mu\nu}^\alpha = \partial_\mu A_\nu^\alpha - \partial_\nu A_\mu^\alpha - g_1 \epsilon^{\alpha\beta\gamma} A_\beta^\gamma A_\mu^\alpha, \]  
(38)

\[ B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu, \]  
(39)

for the three non-Abelian fields of \( SU(2)_L \) and the single Abelian gauge field associated with \( U(1)_Y \), respectively. The covariant derivative is:
\[ D_\mu = \partial_\mu - \frac{1}{2} ig_1 A_\mu^a T_a + \frac{1}{2} ig_1^{'} B_\mu, \]  
(40)

with \( g_1, g_1^{'} \) being the \( SU(2)_L, \) and \( U(1)_Y \) the coupling strength, respectively. This Lagrangian is invariant under the infinitesimal local gauge transformations for \( SU(2)_L \) and \( U(1)_Y \) independently. Being in the adjoined representation, the \( SU(2)_L \) massless gauge fields form a weak isospin triplet, with the charged fields being defined by
\[ W_\mu^\pm = (A_\mu^1 \mp i A_\mu^2)/\sqrt{2}. \]  
(41)

The neutral component of \( A_3^\mu \) mixes with the Abelian gauge field \( B_\mu \) to form the physical states:
\[ Z_\mu = A_3^\mu \cos \theta_w - B_\mu \sin \theta_w, \]  
(42)

\[ A_\mu = B_\mu \cos \theta_w + A_3^\mu \sin \theta_w, \]  
(43)

where \( \tan \theta_w = \frac{g_1^{'}}{g_1} \) is the weak mixing angle.

Based on the gauge group \( SU(N) \times U(2) \), the final Lagrangian of the model is given as follows:
\[ S_{Model} = \bar{\psi} \gamma^\mu \left[ \partial_\mu - i g_1 \frac{1}{2} \tau \cdot A_\mu + i g'_1 \frac{1}{2} B_\mu - i g_2 a_\mu^I T_I \right] \psi + \bar{e}_R \gamma^\mu \left[ \partial_\mu + i g'_1 B_\mu \right] e_R \]
\[
- \frac{1}{4} F^a_{\mu \nu} F^{a \mu \nu} - \frac{1}{4} B_{\mu \nu} B^{\mu \nu} - \frac{1}{4} a^{I \mu \nu} a^I_\mu - \frac{1}{4} b^{I \mu \nu} b^I_\mu
\]
\[ - \frac{\mu^2}{2(1 + c^2)} (a^I_\mu + cb^I_\mu)(a^I_\mu + cb^I_\mu) \quad (44) \]

where \(c\) is a constant.

3 The Masses of Gauge Fields

Two obvious characteristics of the Wu Lagrangian equation (36) is that the mass term of the gauge fields is introduced into the Lagrangian, and that this term does not affect the symmetry of the Lagrangian. It has been proved that this Lagrangian has strict local gauge symmetry [16]. Since both vector fields \(a_\mu\) and \(b_\mu\) are standard gauge fields, this model is a gauge field model which describes gauge interactions between gauge fields and matter fields [16].

The mass term of gauge fields can be written as follows [16]:
\[ \frac{1}{2} (a^\mu, b^\mu) M \begin{pmatrix} a^\mu \\ b^\mu \end{pmatrix}, \quad (45) \]
where \(M\) is the mass matrix:
\[ M = \frac{1}{1 + c^2} \begin{pmatrix} \mu^2 & c \mu^2 \\ c \mu^2 & c^2 \mu^2 \end{pmatrix}. \quad (46) \]

Physical particles generated from gauge interactions are eigenvectors of mass matrix, and the corresponding masses of these particles are eigenvalues of mass matrix. The mass matrix \(M\) has two eigenvalues:
\[ m_1^2 = \mu^2, \quad m_2^2 = 0. \quad (47) \]
The corresponding eigenvectors are:
\[ \begin{pmatrix} \cos \theta_wu \\ \sin \theta_wu \end{pmatrix} \begin{pmatrix} -\sin \theta_wu \\ \cos \theta_wu \end{pmatrix} \quad (48) \]
where
\[ \cos \theta_wu = \frac{1}{\sqrt{1 + c^2}}, \quad \sin \theta_wu = \frac{c}{\sqrt{1 + c^2}}. \quad (49) \]

We define
\[ C_\mu = \cos \theta_wu a_\mu + \sin \theta_wu b_\mu, \quad (50) \]
\[ F_\mu = -\sin \theta_wu a_\mu + \cos \theta_wu b_\mu. \]

\(C_\mu\) and \(F_\mu\) are eigenstates of mass matrix: they describe the particles generated from gauge interactions. The inverse transformations of (50) are:
\[ a_\mu = \cos \theta_wu C_\mu - \sin \theta_wu F_\mu, \quad (51) \]
\[ b_\mu = \sin \theta_wu C_\mu + \cos \theta_wu F_\mu. \]
Taking (50) and (51) into account, the Wu Lagrangian density $L$ given by (36) changes into:

$$\mathcal{L}_{\text{Wu}} = \mathcal{L}_{\text{Wu}}^{(0)} + \mathcal{L}_{\text{Wu}}^{(I)},$$

where

$$\mathcal{L}_{\text{Wu}}^{(0)} = -\bar{\psi}(\gamma^\mu \partial_\mu + m)\psi - \frac{1}{4} C_{\mu \nu}^i C_{i \mu \nu}^j - \frac{1}{4} K_{\mu \nu}^i K_{i \mu \nu}^j - \frac{\mu^2}{2} C_{i \mu}^i C_{\mu}^i,$$

$$\mathcal{L}_{\text{Wu}}^{(I)} = ig_2 \bar{\psi} \gamma^\mu (\cos \theta_{\text{wu}} C_{\mu}^i - \sin \theta_{\text{wu}} F_{\mu}) \psi - \frac{\cos 2\theta_{\text{wu}}}{2 \cos \theta_{\text{wu}}} g_2 f_{ij} f_{k\mu} C_{i \mu}^j C_{k \nu}^j + g_2 \sin \theta_{\text{wu}} f_{ij} C_{i \mu}^j C_{\mu}^k F_{\nu}^k - \frac{1 - \frac{3}{4} \sin^2 2\theta_{\text{wu}}}{4 \cos^2 \theta_{\text{wu}}} g_2^2 f_{ij} f_{k \mu} f_{\nu}^l f_{\mu}^l C_{\nu}^i C_{i \nu}^j C_{\mu}^k F_{\nu}^k - \frac{\sin^2 \theta_{\text{wu}}}{2} g_2^2 f_{ij} f_{k \mu} f_{\nu}^l f_{\nu}^l C_{i \mu}^j C_{\nu}^i C_{\nu}^j C_{\mu}^k F_{\nu}^k.
$$

In the above relations, we have used the following simplified notations:

$$C_{0 \mu \nu}^i = \partial_\mu C_{i \nu}^j - \partial_\nu C_{i \mu}^j,$$

$$K_{0 \mu \nu}^i = \partial_\mu F_{i \nu}^j - \partial_\nu F_{i \mu}^j.$$

From (53) it is deduced that the mass of field $C_{\mu}$ is $\mu$ and the mass of gauge field $F_{\mu}$ is zero. That is:

$$M_C = \mu, \quad M_F = 0,$$

Transformations (50) and (51) are pure algebraic operations which do not affect the gauge symmetry of the Lagrangian [16]. They can, therefore, be regarded as redefinitions of gauge fields. The local gauge symmetry of the Lagrangian is still strictly preserved after field transformations. In other words, the symmetry of the Lagrangian before transformations is absolutely the same with the symmetry of the Lagrangian after transformations. We do not introduce any kind of symmetry breaking at energy scales close to 2 TeV.

Fields $C_{\mu}$ and $F_{\mu}$ are linear combinations of gauge fields $a_{\mu}$ and $b_{\mu}$. The forms of local gauge transformations of fields $C_{\mu}$ and $F_{\mu}$ are, therefore, determined by the forms of local gauge transformations of gauge fields $a_{\mu}$ and $b_{\mu}$. Because $C_{\mu}$ and $F_{\mu}$ consist of gauge fields $a_{\mu}$ and $b_{\mu}$ and transmit gauge interactions between matter fields, for the sake of simplicity we also call them gauge fields, just as $W$ and $Z$ are called gauge fields in the electroweak model [20–23]. This gauge field theory, therefore, predicts the existence of two different kinds of force transmitting vector fields: of a massive and a massless one.

Taking the Higgs mechanism [20–23] into account, in the vacuum energy scale of 179 GeV the $W^\pm$ and $Z^0$ become massive, while the photon $A$ remains massless. The symmetry $SU(N)$ does not break down, since the gauge bosons $C^0$ and $F$ derive their masses by the Wu mechanisms [16–19] in the vacuum energy scale of 2 TeV.

The most general Lagrangian consistent with $SU(2) \times U(1)$ gauge invariance, Lorentz invariance, and with renormalizability is:

$$\mathcal{L}_\phi = -\frac{1}{2} |(\partial_{\mu} - i A_{\mu} \gamma^\phi - i B_{\mu} y(\phi)) \phi|^2 - \frac{\mu^2}{2} \phi^* \phi - \frac{\lambda}{4} (\phi^* \phi)^2$$

([32]).

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Table 1  Gauge fields and masses predicted by the proposed model

| Gauge fields | Masses (M)-GeV | Symmetry pattern               |
|--------------|---------------|--------------------------------|
| C (New)      | 145           | Without symmetry breaking      |
| F (New)      | 0             | Without-symmetry breaking      |
| W            | 81            | With-symmetry breaking         |
| Z            | 91            | With-symmetry breaking         |
| A            | 0             | With-symmetry breaking         |

where

\[
t^{(\phi)} = \frac{g_1}{2} \left\{ \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} \right\}, \quad y^{(\phi)} = -\frac{g_1'}{2} \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}
\]

(59)
generators of the \((\phi^+, \phi^0)\), \(\lambda > 0\), and

\[
\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}.
\]

(60)

For \(\mu^2 < 0\), there is a tree-approximation vacuum expectation value at the stationary point of the Lagrangian

\[
\langle \phi^\ast \rangle \langle \phi \rangle = v^2 = |\mu^2|/\lambda.
\]

(61)

We can always perform an \(SU(2) \times U(1)\) gauge transformation to a unitary gauge, in which \(\phi^+ = 0\) and \(\phi^0\) is Hermitian, with positive vacuum expectation value. In unitarily gauge the vacuum expectation values of the components of \(\phi\) are

\[
\langle \phi^+ \rangle = 0, \quad \langle \phi^0 \rangle = v > 0.
\]

(62)

The scalar Lagrangian (58) then yields a vector boson mass term:

\[
-\frac{1}{2} \left| \left( \partial \mu - i A_\mu \cdot \tau^{(\phi)} - i B_\mu y^{(\phi)} \right) \phi \right|^2 = -\frac{1}{2} \left| \left( \frac{g_1}{2} A_\mu T^a - \frac{g_1'}{2} B_\mu \right) \begin{pmatrix} 0 \\ v \end{pmatrix} \right|^2 = -\frac{v^2 g_1^2}{4} W_\mu W^\mu - \frac{v^2}{8} (g_1^2 + g_1'^2) Z_\mu Z^\mu.
\]

(63)

The masses are given as follows

\[
M_A = 0, \quad M_W = \frac{1}{2} g_1 v, \quad M_Z = \frac{1}{2} \sqrt{g_1^2 + g_1'^2}.
\]

(64)

The gauge fields and masses predicted by this model are summarized in Table 1.

4 The Renormalization of the Model

In this paper, we use two mechanisms that can make gauge field to gain nonzero mass. One is the Wu mechanism [16–19, 30, 31]. In this application of the Wu mechanism, by which the mass term of the gauge field is introduced by using another set of gauge field. In this mechanism, the mass term of the gauge field does not affect the symmetry of the Lagrangian. We can imagine the new interaction picture: when matter fields take part in gauge interactions, they emit or absorb one kind of gauge field which is not eigenstate of mass matrix. Perhaps this is the gauge field detected in the CDF experiments [2]. This gauge
field would appear in two states, a massless and a massive one, which correspond to two kinds of vector fields.

The other is the Higgs mechanism [20–23], which can make gauge field—\( W, Z \) to gain nonzero mass and to guarantee renormalizability by means of the interactions of the Higgs boson with gauge bosons:

\[
\frac{\eta^2 + 2v\eta}{v^2}(Z_\mu Z^\mu M^2_Z + C_\mu C^\mu M^2_C + W^+_\mu W^{\mu-} M^2_W) \\
= \left( \frac{M^2_Z}{v^2} \right) \eta^2 Z_\mu Z^\mu + \left( \frac{M^2_C}{v^2} \right) \eta^2 C_\mu C^\mu \\
+ \left( \frac{M^2_W}{v^2} \right) \eta^2 W^+_\mu W^{\mu-} + 2 \left( \frac{M^2_Z}{v} \right) \eta Z_\mu Z^\mu \\
+ 2 \left( \frac{M^2_C}{v} \right) \eta C_\mu C^\mu + 2 \left( \frac{M^2_W}{v} \right) \eta W^+_\mu W^{\mu-}
\]

(65)

where

\[
M^2_Z/v^2, M^2_C/v^2, M^2_W/v^2, 2M^2_Z/v, 2M^2_C/v, 2M^2_W/v
\]

(66)

the dimensionless and dimension of mass coupling constants.

For instance, the \( C \)-boson readily derives its mass by the Wu mechanism and yet the renormalizability is ensured via the Higgs mechanism.

However, as we have stated above, the Wu gauge field theory has maximum local \( SU(N) \) gauge symmetry [16]. When we quantize the Wu gauge field theory in the path integral formulation, we must select gauge conditions first [16]. To fix the degree of freedom of the gauge transformation, we must select two gauge conditions simultaneously: one for the massive gauge field \( C_\mu \) and another for the massless gauge field \( F_\mu \). For example, if we select temporal gauge condition for massless gauge field \( F_\mu \),

\[
F_4 = 0,
\]

(67)

there still exists a remainder gauge transformation degree of freedom, because the temporal gauge condition is unchanged under the following local gauge transformation:

\[
F_4 \rightarrow U F_\mu U^+ + (1/ig_2 \sin \theta) U \partial_\mu U^+.
\]

(68)

where

\[
\partial_t U = 0, \quad U = U(\vec{x}).
\]

(69)

In order to make this remainder gauge transformation degree of freedom completely fixed, we have to select another gauge condition for gauge field \( C_\mu \). For instance, we can select the following gauge condition for gauge field \( C_\mu \),

\[
\partial^\mu C_\mu = 0.
\]

(70)

If we select two gauge conditions simultaneously, when we quantize the theory in path integral formulation, there will be two gauge fixing terms in the effective Lagrangian. The effective Lagrangian can be written as:

\[
\mathcal{L}_{\text{eff}} = \mathcal{L} - \frac{1}{2\alpha_1} f_1^a f_1^a - \frac{1}{2\alpha_2} f_2^a f_2^a + \tilde{\eta}_1 M_{f_1} \eta_1 + \tilde{\eta}_2 M_{f_2} \eta_2.
\]

(71)

where

\[
f_1^a = f_1^a (F_\mu), \quad f_2^a = f_2^a (C_\mu).
\]

(72)
If we select
\[ f_a^a = \partial^\mu C^{a\mu}_\mu, \quad (73) \]
then the propagator for massive gauge field \( C^\mu \) is:
\[
\Delta_{F^{ab}}^{\mu\nu}(k) = -i\delta^{ab}/(k^2 + \mu^2 - i\varepsilon) \times \left[ g_{\mu\nu} - (1 - \alpha_2)k_\mu k_\nu/(k^2 + \alpha_2\mu^2) \right]. \quad (74)
\]
If we let \( k \) approach infinity, then
\[
\Delta_{F^{ab}}^{\mu\nu}(k) = \frac{1}{k^2}. \quad (75)
\]
In this case, according to the power-counting law, the Wu gauge field theory suggested in this paper is a kind of renormalizable theory [16].

5 The Baryonic C-Model

In this paper, we propose a baryonic C model to explain the Wij anomaly. The reason this model to be baryonic is that, even if \( C \) has a small leptonic branching ratio (even \( O(1) \% \)), it would suffer from the strong constraints of the Tevatron \( C \) search in the dilepton mode [40]. The proposed \( C \) gauge boson is similar to the \( Z \)-prime, a hypothetical carrier of a new force similar to the electroweak force [29].

The baryonic \( Z' \) model was proposed by Barger, Cheung, and Langacker in 1996 [41] in light of the \( R_b/R_c \) crisis of the LEP precision measurements at that time [42, 43]. That interpretation of \( C \) at the time suggests strong implications for the Tevatron [41] via \( s \)-channel \( C \) production and the pair production processes \((W, Z, \gamma)C\), with \( C \rightarrow jj \) (in particular \( b\bar{b} \)), with invariant mass \( M_{jj} \) peaked at \( M_C \). The \( s \)-channel \( C \) production is buried under the QCD background, but the associated production with a \( W \) boson has a good chance to appear.

The current CDF anomaly [2] may be of this origin. Here, as a consequence of the baryonic \( N \)-component vector in the fundamental representation of \( SU(N) \) group, the gauge field \( b^\mu_i \) is also baryonic. Thus, by \((50)\), the gauge fields \( F \) and \( C \) are baryonic. The four physical neutral gauge fields \( A^{\mu}_1, Z^{\mu}_1, F^{\mu}_1, C^{\mu}_1 \) are orthogonal combinations of the gauge fields \((A^3_3, B_3)\) and \((a^\mu_2, b^\mu_2)\) with mixing angles \( \theta_w, \theta_{wu} \) respectively.

The Lagrangian describing the neutral current gauge interactions of the gauge group \( SU(N) \times U(2) \) (Eq. \((1)\)), in terms of these physical fields is given as follows:
\[
\mathcal{L}_{NC}^{(Model)} = \mathcal{L}_{NC}^{(Wu)} + \mathcal{L}_{NC}^{(SM)}, \quad (76)
\]
where
\[
\mathcal{L}_{NC}^{(SM)} = -ie\left( j^{em}\right)^\mu A_\mu - ig_1 Z^{\mu}_0 \left[ \frac{1}{2} \bar{\psi} \gamma^\mu (g^{(1)}_\nu - g^{(1)}_a \gamma^5) \psi \right], \quad (77)
\]
\[
\mathcal{L}_{NC}^{(Wu)} = -ib^\mu F_\mu - ig_2 C^{\mu}_0 \left[ \frac{1}{2} \bar{\psi} \gamma^\mu (g^{(2)}_\nu - g^{(2)}_a \gamma^5) \psi \right]. \quad (78)
\]
The electroweak and Wu neutral currents, \( b \) is the baryonic number and \( e \) the electric charge. \( g_1 \) and \( g_2 \) are the SM and the \( C \) boson coupling constants, respectively. Substituting \((77)\) and \((78)\) to \((76)\) we get:
\[
-\mathcal{L}_{NC}^{(Model)} = e\left( j^{em}\right)^\mu A_\mu + bj^\mu F_\mu + g_1 Z^{\mu}_0 \left[ \frac{1}{2} \bar{\psi} \gamma^\mu (g^{(1)}_\nu - g^{(1)}_a \gamma^5) \psi \right]
\]
\[
+ g_2 C^{\mu}_0 \left[ \frac{1}{2} \bar{\psi} \gamma^\mu (g^{(2)}_\nu - g^{(2)}_a \gamma^5) \psi \right]. \quad (79)
\]
The coupling constant $g_1$ is the SM coupling $g / \cos \theta_w$. For grand unified theories (GUT) $g_2$ is related to $g_1$ by

$$
\frac{g_2}{g_1} = \left( \frac{5}{3} x_w \lambda \right)^{1/2} \approx 0.62 \lambda^{1/2} \quad (80),
$$

where $x_w = \sin^2 \theta_w$ and $\theta_w$ is the weak mixing angle. The factor $\lambda$ depends on the symmetry breaking pattern and the fermion sector of the theory, which is usually of order unity. Since we only consider the mixing of $Z^0$ and $C^0$, we can rewrite the Lagrangian with only the $Z^0$ and $C^0$ interactions:

$$
-V_{ZC}^{(\text{Model})} = g_1 Z_{\mu} \left[ \frac{1}{2} \bar{\psi} \gamma^\mu \left( g^{(1)}_u - g^{(1)}_a \gamma^5 \right) \psi \right] + g_2 C_{\mu} \left[ \frac{1}{2} \bar{\psi} \gamma^\mu \left( g^{(2)}_u - g^{(2)}_a \gamma^5 \right) \psi \right],
$$

where for both quarks and leptons it is

$$
g^{(1)}_u = T_{3L} - 2 x_w Q, \quad g^{(1)}_a = T_{3L}. \quad (82)
$$

We consider the case in which $C^0$ couples only to quarks,

$$
g^{(2)}_u = \epsilon_V = \sin \gamma, \quad g^{(2)}_a = \epsilon_A = \cos \gamma, \quad g^{(2)}_a = g^{(2)}_a = 0. \quad (83)
$$

Here $T_{3L}$ is the third component of the weak isospin, and $Q$ is the electric charge of the fermions. The parameters $\epsilon_V$ and $\epsilon_A$ are the vector and axial-vector couplings of $C^0$. Without loss of generality we choose $\epsilon_V = \sin \gamma$ and $\epsilon_A = \cos \gamma$, such that $(\epsilon_V^2 + \epsilon_A^2)$ is normalized to unity. The mixing of the weak eigenstates $Z^0$ and $C^0$ to form mass eigenstates $Z$ and $C$ is parametrized by a mixing angle $\theta$:

$$
Z = Z^0 \cos \theta + C^0 \sin \theta, \quad C = C^0 \cos \theta - Z^0 \sin \theta. \quad (84)
$$

After substituting the interactions of the mass eigenstates and $Z'$ with fermions we obtain

$$
-V_{ZC}^{(\text{Model})} = \frac{g_1}{2} \left[ Z_{\mu} \bar{\psi} \gamma^\mu \left( u_s - a_s \gamma^5 \right) \psi + C_{\mu} \bar{\psi} \gamma^\mu \left( \nu_a - a_n \gamma^5 \right) \psi \right],
$$

where

$$
\begin{align*}
\nu_s &= g^{(1)}_u + \frac{g_2}{g_1} \theta g^{(2)}_u, \\
\alpha_s &= g^{(1)}_a + \frac{g_2}{g_1} \theta g^{(2)}_a,
\end{align*}
$$

$$
\begin{align*}
\nu_n &= \frac{g_2}{g_1} g^{(2)}_u - \theta g^{(1)}_u, \\
\alpha_n &= \frac{g_2}{g_1} g^{(2)}_a - \theta g^{(1)}_a.
\end{align*}
$$

Here we have used the valid approximation $\cos \theta \approx 1$ and $\sin \theta \approx \theta$.

In the following, we ignore the mixing ($\theta = 0$), so that the precision measurements for the SM $Z$ boson are not affected, unless stated otherwise. We also take the democratic choice of equal couplings of $C$ to up-type and down-type quarks. This is in accord with the CDF observation that there is no preference for $b$ quarks in the dijet window $M_{jj} = 120–160$ GeV

$$
\Gamma(C \to f \bar{f}) = \frac{G_F M_C^2}{6 \pi \sqrt{2}} N_c C \left( M_C^2 \right) M_C \sqrt{1 - 4 \chi} \left[ u_n^2 (1 + 2 \chi) + a_n^2 (1 - 4 \chi) \right], \quad (88)
$$

where $G_F$ is the Fermi coupling constant, and $N_c = 3$ or 1, if $f$ is a quark or a lepton, respectively.

The decay width of the $C$ gauge boson into fermions is given by

$$
\Gamma(C \to f \bar{f}) = \frac{G_F M_C^2}{6 \pi \sqrt{2}} N_c C \left( M_C^2 \right) M_C \sqrt{1 - 4 \chi} \left[ u_n^2 (1 + 2 \chi) + a_n^2 (1 - 4 \chi) \right],
$$

$$
C \left( M_C^2 \right) = \alpha_s / \pi + 1.409 (\alpha_s / \pi)^2 - 12.77 (\alpha_s / \pi)^3, \quad (89)
$$

$$
\alpha_s = \alpha_s (M_C).
$$
The term $C$ represents the strong coupling at the scale $M_C$, $x = m_f^2/M_C^2$. The $C$ width is proportional to $\lambda$, which sets the strength of the $C$ coupling. For $\lambda = 1$ the total $C$ width is
\[ \Gamma_C/M_C = 0.022 \quad \text{for } M_C < 2m_t. \quad (90) \]

This width is increased somewhat if there are open channels for decay into the top quark, superpartners, and other exotic particles.

The $C$ boson can be produced directly at a hadron collider via the $\bar{q}q \rightarrow C$ sub-process, for which the cross section in the narrow $C$ width approximation is:
\[ \hat{\sigma}(\bar{q}q \rightarrow C) = K \frac{2\pi}{3} \frac{G_F M_C^2}{\sqrt{2}} \left[ (u_q^2) + (a_q^2) \right] \delta(\hat{s} - M_C^2). \quad (91) \]

The $K$-factor represents the enhancement from higher-order QCD processes, estimated to be:
\[ K = 1 + \frac{\alpha_s(M_C^2)}{2\pi} \left( 1 + \frac{4}{3}\pi^2 \right) \approx 1.3 \quad ([24]). \quad (92) \]

If mixing is ignored it is,
\[ (u_q^2) + (a_q^2) = (0.62)^2 \lambda. \quad (93) \]

The above cross section is independent of the parameter $\gamma$ as long as $\epsilon_V^2 + \epsilon_A^2 = 1$.

Note that all current and previous dijet-mass searches at the Tevatron are limited to $M_{jj} > 200$ GeV. These energy levels are not applicable to the present $C$ with $M_C \approx 145$ GeV. The relevant dijet data are those from the UA2 Collaboration, with collision energy at $\sqrt{s} = 630$ GeV [25]. The UA2 Collaboration has detected the $W + Z$ signal in the dijet mass range $48 < m(jj) < 138$ GeV, and has placed the upper boundary of $\sigma(B(C \rightarrow jj))$ over the range $80 < m(jj) < 320$ GeV.

The associated production of $C$ with a $W$-boson goes through the $t$- and $u$-channel exchange of quarks, while the s-channel boson exchange is highly suppressed because of the negligible mixing angle between the SM $Z$-boson and the $C$. Consequently, we expect similar, or even larger cross sections for $M_C \sim M_Z$ than the SM $WZ$ production in which there is a delicate gauge cancellation among the $t$-, $u$-, and $s$-channel diagrams. We have included a $K$-factor ($K = 1.3$) to approximate next-to-leading order QCD contributions [26]. We can see that, at around 140–150 GeV, the cross section is right at the order of 4 pb, which is required to explain the excess in the CDF $Wjj$ anomaly [2].

### 6 Conclusion

We have shown that a new neutral vector $C$-boson of mass (145 GeV), predicted by the Wu mechanisms for mass generation of gauge field, can explain the excess in the invariant-mass window 120–160 GeV in the dijet system of $Wjj$ production.

In this model, the Standard Model (SM) $W, Z$-bosons acquire their masses through the coupling with the SM Higgs of mass 114–200 GeV. The $C$-boson has negligible couplings to leptons, and so is not affected by the dilepton $C$ constraints.

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