Production of new chiral bosons at Tevatron and LHC

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The possibility of production of new spin-1 chiral particles at the hadron colliders the Tevatron and the LHC is considered. Their coupling constants to ordinary fermions can be fixed using the hypothesis about a dynamical generation of kinetic terms for the bosons and assuming an universality of interactions of these particles. The masses of the chiral particles can be predicted from experimental data of precise low-energy experiments, which already indicate indirectly the existence of such new particles. Quantitative estimations of the production cross-sections of the chiral particles are made. It is noteworthy that, due to their non-gauge interactions with fermions, their decays do not lead to the Jacobian peak in the transverse momentum/mass distribution. Moreover, the lightest chiral particles show leptophobic character, which makes their detection at hadron colliders challenging.

I. INTRODUCTION

The standard model (SM) contains and unifies interactions of three different types of elementary particles, with respect to spin: scalars, spin one-half fermions and vector gauge bosons. The quantum field theory for such types of particles has been successfully constructed and has passed through severe experimental examinations. Almost all phenomenological attempts to extend the SM consist in introducing of a number of new particles of the same type, like extra Higgs multiplets, fourth fermion generation and various $Z'$ gauge bosons. Even SUSY phenomenology does not go further simple doubling of known types of particles.

In this letter we would like to discuss again the possibility of the presence in nature of new types of spin-1 chiral bosons, which are complementary to the gauge ones. While the latter transform according to the real representation $(1/2,1/2)$ of the Lorentz group, transformation properties of the former ones are associated with the inequivalent chiral representations $(1,0)$ and $(0,1)$. The chiral spin-1 bosons are described by rank-two antisymmetric tensor fields $T_{mn}$, instead of vector potential $V_m$, which describes the gauge spin-1 bosons.

On the mass shell both fields describe spin-1 particles and an equivalence relation between them can be set in the massive case

$$T_{mn} = \frac{1}{M} (\partial_m V_n - \partial_n V_m) - \frac{\epsilon_{mnab}}{2M} (\partial_a A_b - \partial_b A_a),$$

where $M$ is the mass of spin-1 particles. The tensor field describes simultaneously two three-component vector $V_m$ and axial-vector $A_m$ fields. They can be expressed through six-component tensor field $T_{mn}$, using the on-shell operator equation $\partial^2 = -M^2$,

$$V_m = \frac{1}{M} \partial_m T_{mn}, \quad A_m = \frac{1}{2M} \epsilon_{mnab} \partial_n T_{ab}. \quad (2)$$

However, many years ago it was noted that different Yukawa interactions for different types of spin-1 bosons lead to inequivalent models. The first successful phenomenological application of this idea has been realized for hadron meson states with spin one [2].

It was pointed out that three different quantum numbers $J^{PC}$ of existing spin-1 mesons, $1^{-+}$, $1^{++}$ and $1^{+-}$, cannot be assigned just to two vector $q\gamma_m q$ and axial-vector $\vec{q}\gamma_m \gamma^5 q$ quark states. So, the rank-two anti-symmetric tensor current $\bar{q}\sigma_{mn} q$ was introduced, which also describes vector and axial-vector boson states, but with different transformation properties with respect to Lorentz group and with different quantum numbers $1^{-+}$ and $1^{+-}$, respectively. This example demonstrates that both the pure tensor states, $b_1$ bosons, and mixed combinations of vector and tensor states, due to the same quantum numbers, $\rho$ and $\rho'$ bosons, may exist.

These considerations, in general, support the idea of extension of the SM with a new type of spin-1 chiral particles [3], which has been proposed on pure theoretical grounds, in order to describe the anomalies in the weak charged meson decays. Assuming $SU_L(2) \times SU_Y(1)$ symmetry of electroweak interactions and that very heavy right-handed neutrinos $\nu_{Ra}$ are decoupled at the weak scale, most general Yukawa interactions were proposed

$$L_Y = \frac{1}{4} \left[ t^2_a (\bar{Q}_a \sigma^{mn} d_{Ra}) + t'_a (\bar{L}_a \sigma^{mn} \epsilon_{Ra}) \right] \left( \frac{T_{mn}^+}{T_{mn}^0} \right)$$

$$+ \frac{u^2}{4} (\bar{Q}_a \sigma^{mn} u_{Ra}) \left( \frac{u_{mn}}{U_{mn}} \right) + h.c., \quad (3)$$

where the two doublets $T_{mn}$ and $U_{mn}$ of the new tensor chiral fields with opposite hypercharges $Y(T) = +1$, $Y(U) = -1$ were introduced.

The fermion sector is the same as in the SM and consists of three generations (noted by the index $a$) of the two component Weyl spinors $\psi_a$: the left-handed lepton $L_a$ and the color quark $Q_a$ doublets and the right-handed lepton $\epsilon_{Ra}$ and the color quark $u_{Ra}$, $d_{Ra}$ singlets. Here and further on we omit quark color indexes as far as the tensor fields are color blind. The local gauge electroweak symmetry, maintained by the triplet $\mathbf{W}_m$ and
singlet $B_m$ bosons, and $SU_C(3)$ color symmetry rest untouched. It was also argued \cite{5} that in order to avoid new chiral anomalies connected with the antisymmetric chiral tensor fields, two Higgs doublets with opposite hypercharges should be introduced instead of the one in the SM.

Topologically the interactions \cite{25} are the same as the electroweak gauge interactions

$$L_{\text{gauge}} = g \left( \bar{\psi}_a \gamma^m T \psi_a \right) W_m + g' \left( \bar{\psi}_a \gamma^m \frac{Y}{2} \psi_a \right) B_m, \quad (4)$$

where $g$, $g'$ are the gauge coupling constants, and $T$, $Y$ are the generators of $SU_L(2)$ and $U_Y(1)$, correspondingly. The only principal difference between them concerns fermion chiralities. While the gauge field interactions preserve the helicities of incoming and outcoming fermions, the tensor field couplings lead to a helicity flip. This feature should be the main signature in their experimental search.

II. COUPLING CONSTANT RELATIONS

In this section we will obtain relations among the gauge coupling constants $g$, $g'$ and the tensor Yukawa couplings $t^u_\alpha$, $t^d_\alpha$ and $u^\alpha$. In order to do this we will apply hypothesis about dynamical generation of kinetic terms for the bosons. This idea comes from Nambu and Jona-Lasinio model \cite{5} with nonlinear four-fermion interactions. It has been developed in \cite{5} for the more appropriate for quantization linear form of interactions, introducing auxiliary boson fields (without kinetic terms) with dimensionless Yukawa coupling constants $\hat{v}^u$, $\hat{v}^d$.

Let us first apply this method for the gauge fields $W_m$ and $B_m$, assuming that they do not posses initial kinetic terms. The kinetic terms are generated by self-energy quantum corrections from fermionic loops. Subsequent renormalization procedure leads to proper normalization of the kinetic terms and sets relations among various Yukawa coupling constants.

So, in the lowest one-loop approximation the self-energy quantum corrections to the gauge fields $W_m$ read

$$\Pi_{W_{mn}}(q) = ig^2(1 + N_C)N_g \text{Tr} \left( T^i T^j \right) \times \int \frac{d^4 p}{(2\pi)^4} \text{Tr} \left[ \gamma_m \frac{1 - \gamma^5}{2} (\hat{p} - \hat{q})^{-1} \gamma_n \frac{1 - \gamma^5}{2} (\hat{p} - \hat{q})^{-1} \right]$$

$$= -4g^2 I_0 \left( g_{mn}q^2 - q_m q_n \right) \delta^{ij} + \mathcal{O}(q^4), \quad (5)$$

where $I_0 = -i/(2\pi)^4 \text{Reg} \int d^4 p/p^4$ is the regularized value of the logarithmically divergent integral. Analogous calculations for the abelian gauge field $B_m$ give

$$\Pi_{B_{mn}}(q) = -\frac{20}{3} g'^2 I_0 \left( g_{mn}q^2 - q_m q_n \right) + \mathcal{O}(q^4). \quad (6)$$

The comparison of the two previous equations leads to a relation between the gauge coupling constants

$$\left( \frac{g'}{g} \right)^2 = \frac{3}{5} \quad (7)$$

and, hence, to a prediction of the electroweak mixing angle $\sin^2 \theta_W = 0.23152(14)$. A lack of a gauge symmetry for the chiral tensor interactions \cite{25} does not allow to fix the Yukawa coupling constants $t^u_\alpha$, $t^d_\alpha$ and $u^\alpha$. Therefore, to go further we must make model dependent assumptions. Let us assume lepton-quark universality for the tensor interactions, which are also blind to the generations $t^\ell_\alpha = t_\alpha$, $u^\alpha = u$.

Before proceeding with calculations of self-energy quantum corrections to the tensor fields, it is convenient to rewrite them through the more common notations of the vector potential. This is possible using extension of the relation \cite{5} off-shell \cite{2}.

$$T_{mn} = \left( \delta_{mN}V_n - \delta_{nN}V_m \right) - \frac{\epsilon_{mnab}}{2} \left( \delta_{aN}A_b - \delta_b A_a \right), \quad (9)$$

where $\delta_m = \partial_m/\sqrt{-\hat{D}^2}$. Then eq. \cite{5} takes the form

$$L_Y = t \left( \bar{Q}_a \sigma^{mn} d_{Ra} + \bar{L}_a \sigma^{mn} e_{Ra} \right) \left( \delta_{mN} T^+_{nN} - \delta_{nN} T^+_{mN} \right) + u \left( \bar{Q}_a \sigma^{mn} u_{Ra} \right) \left( \delta_{mN} U^0_{nN} - \delta_{nN} U^0_{mN} \right) + \text{h.c.}, \quad (10)$$

Now it is easy to calculate the self-energy quantum corrections to doubles $T_m$

$$\Pi_{T_{mn}}(q) = -4t^2 I_0 \left( g_{mn}q^2 - q_m q_n \right) + \mathcal{O}(q^4) \quad (11)$$

and $U_m$

$$\Pi_{U_{mn}}(q) = -3u^2 I_0 \left( g_{mn}q^2 - q_m q_n \right) + \mathcal{O}(q^4). \quad (12)$$

The comparison of \cite{10}, \cite{11} and \cite{12} leads to the relation

$$g^2 = t^2 = \frac{3}{4} u^2. \quad (13)$$

An asymmetry in the coupling constants for the doubles of the tensor fields arises due to their obviously non-symmetric interactions to the leptons. At high energy scale, when the right-handed neutrinos become active, the symmetry between both interactions should be restored.

It may be expected that the evolution of the coupling constants towards low-energy scale should not drastically change the relation \cite{13}, since both the gauge coupling constant $g$ and the Yukawa coupling constants $t$ and $u$ have asymptotically free behavior \cite{10}. We will use this relation below for a quantitative determination of the masses of the chiral particles.
III. MASSES AND MIXINGS

The Fermi theory of the weak interaction \[11\] gives excellent example of effective field theory which perfectly describes low-energy physics with only one dimensional coupling constant \(G_F = 1.16637(1) \times 10^{-5} \text{ GeV}^{-2}\). It arises from the exchange of the charged massive intermediate bosons \(W^\pm\) at small momentum transfer, which at the tree level is expressed as

\[
\frac{G_F}{\sqrt{2}} = \frac{g^2}{8M_W^2},
\]

where \(M_W\) is the boson mass. In the SM the gauge coupling constant \(g^2 = 4\pi\alpha/\sin^2\theta_W\) is related to the electromagnetic fine-structure constant \(\alpha\) and \(\sin^2\theta_W\). It allows to estimate \(M_W \approx 77.5 \text{ GeV}\) at the tree level with relatively good precision.

Therefore, in order to get a constraint on the mass of the new intermediate boson one can analyze the precision low-energy experiments, looking for an eventual admixture of new forces in the ordinary weak processes. Particularly, the chiral bosons can enhance the rates of the chirally suppressed decay modes.

One may think that \(\pi \rightarrow e\nu\) decay is the ideal process to put constraints on the presence of new currents chirally different from the SM ones. However, for the tensor current the hadron matrix element \(\langle 0 | \bar{q} \sigma_{mn} q | \pi \rangle\) is zero by kinematic reasons, and this decay can restrict only the presence of pseudoscalar currents \[12\], but not tensor ones. Nevertheless, the quark tensor currents can contribute to the rare three-body pion decay \(\pi \rightarrow e\nu\gamma\), where huge experimental anomalies have been detected \[13\].

In order to explain these anomalies, the following effective tensor interactions have been derived \[3\]

\[
\mathcal{L}_{\text{eff}} = -\sqrt{2} f_T G_F \bar{u} \sigma_{ml} d_L \frac{4q^2}{q^2} \bar{e} \sigma^{mn} \nu_L
- \sqrt{2} f'_T G_F \bar{u} \sigma_{ml} d_R \frac{4q^2}{q^2} \bar{e} \sigma^{mn} \nu_L + \text{h.c.},
\]

where \(q\) is the momentum transfer between quark and lepton pairs and \(f_T, f'_T\) are dimensionless coupling constants, which determine the strength of the new tensor interactions relative to the ordinary weak interactions.

These interactions arise from the exchange of the charged chiral \(T^\pm\) and \(U^\pm\) bosons at low energies. The first term in \[15\] is generated due to the non-diagonal elements of mass mixing matrix \[3\]

\[
M^2 = \begin{pmatrix} T^+_m & U^+_m \\ -\mu^2 & m^2 \end{pmatrix} \begin{pmatrix} M^2 & -\mu^2 \\ -\mu^2 & m^2 \end{pmatrix}.
\]

It can be rewritten identically in a more common form without momentum dependence

\[
\bar{u} \sigma_{ml} d_L \frac{4q^2}{q^2} \bar{\nu} \sigma^{mn} \nu_L \equiv \bar{u} \sigma_{mn} d_L \bar{\nu} \sigma^{mn} \nu_L.
\]

This local tensor interaction has been considered early in \(\beta\) decay phenomenology on the same footing as other local four-fermion interactions. However, to my knowledge, the second non-local term in \[16\] has never been used before 1993. Owing to the identity \(\bar{u} \sigma_{mn} d_R \bar{\nu} \sigma^{mn} \nu_L \equiv 0\) it is impossible to construct local tensor interaction with an opposite to \[17\] chirality of quark current.

It is worth noting, that the two terms of the tensor interactions \[15\] arise as a result of the chiral boson exchanges and should be considered simultaneously. The requirement of positivity of the determinant \(\Delta = M^2 m^2 - \mu^4 > 0\) of the square mass matrix \[10\] leads to the positive effective coupling constants

\[
f_T = \frac{g^2 m^2}{4\sqrt{2} G_F \Delta} > 0, \quad f'_T = \frac{g^2 \mu^2}{2 \sqrt{6} G_F \Delta} > 0
\]

and the pseudotensor quark currents \(\bar{u} \sigma^{mn} \nu_L d\) can completely cancel out in \[15\] in the case \(f_T = f'_T\). This cancellation helps to avoid another constraint on the tensor interaction \[17\] following from the pion decay due to generation of the pseudoscalar quark current \(\bar{u} \gamma^5 d\) through the electromagnetic radiative corrections \[14\].

In the case of equality of the effective tensor coupling constants, the diagonalization of the mass matrix \[10\] gives two mass states

\[
M^2_{H/L} = \frac{M^2 + m^2 \pm \sqrt{(M^2 - m^2)^2 + 3m^4}}{2}.
\]

The maximum value of the lightest mass state is reached at \(M^2 = 2m^2\) and defines the physical mass \(M^2_U = m^2/2\), which corresponds to an energetically favored exchange by this particle. Then, the heavy state has a mass \(M^2_H = 5m^2\).

The centi-weak tensor interaction \(f_T \approx 10^{-2}\) can explain the deficit of events in the radiative pion decay and, assuming \(e - \tau\) universality, the excess in \(\tau \rightarrow 2\pi\nu\) decay in comparison with CVC prediction \[17\]. This value of the effective tensor coupling constant allows to predict the masses of the charged chiral bosons

\[
M_H = \sqrt{\frac{2}{f_T}} M_W \approx 1137 \text{ GeV}, \quad M_L \approx 509 \text{ GeV},
\]

which will be used in the next section for quantitative estimations of their production cross-sections.

These physical massive states are represented by two orthogonal combinations \(U^\pm = (\sqrt{3} U^+_n \pm T^+_n)/2\) and \(T^\pm = (\sqrt{3} T^-_n - U^+_n)/2\), which correspond to light and heavy massive particles, correspondingly. Their Yukawa interactions take the form

\[
\mathcal{L}_Y = \frac{g}{2} (\bar{u}_a \sigma^{mn} d_R a + \bar{\nu}_a \sigma^{mn} \nu_L a) \left( \tilde{\partial}_m U^+_n + \sqrt{3} \tilde{\partial}_m T^+_n \right)
+ g (\bar{u}_a \sigma^{mn} d_L a) \left( \tilde{\partial}_m U^+_n - \frac{1}{\sqrt{3}} \tilde{\partial}_m T^+_n \right) + \text{h.c.}.
\]
Besides the mixing \[16\] between the charged chiral particles themselves, they can mix with the \(W^\pm\) boson leading to anomalous moment for the charged weak transitions \[14\]. The first experimental constraints on this parameter have been obtained in \[17\] from \(\tau\)-lepton decays. The mixings between the gauge and the chiral bosons are in a complete analogy with \(\rho - \rho'\) mixing in hadron physics \[2\]. However, in contrast to the latter, where the mixing is maximal, experimental data point out that the weak mixing is rather small.

Let us consider the neutral spin-1 sector of the model. Here besides electroweak bosons \(\gamma\) and \(Z\) we have introduced four degrees of freedom \(T_m^R = (T_m^0 + T_m^\theta)/\sqrt{2}, T_m^L = (T_m^0 - T_m^\theta)/\sqrt{2}, U_m^R = (U_m^0 + U_m^\theta)/\sqrt{2}, U_m^L = (U_m^0 - U_m^\theta)/\sqrt{2}\) for the neutral chiral bosons. Due to their quantum numbers \(T_m^R\) and \(U_m^L\) states do not mix and decouple from the others in the case of \(CP\) invariant interactions. Although \(T_m^R\) and \(U_m^R\) states do not mix between each other, because they couple to the different types up and down fermions, they can mix with the photon and \(Z\) boson leading to additional contributions to anomalous magnetic and neutral weak moments for the fermions. We will assume, in agreement with the experimental data, that these mixings are very small. Therefore, one can consider these states as physical with the following Yukawa interactions

\[
\mathcal{L}_Y = \frac{g}{\sqrt{2}} (\bar{d}_a \sigma^{mn} d_a + \bar{e}_a \sigma^{mn} e_a) \bar{\partial}_m T_m^R + \frac{g}{\sqrt{2}} (\bar{d}_a \sigma^{mn} \gamma^5 d_a + \bar{e}_a \sigma^{mn} \gamma^5 e_a) \bar{\partial}_m T_m^I + \frac{g}{\sqrt{3}} (\bar{d}_a \sigma^{mn} u_a) \bar{\partial}_m U_m^R + \frac{g}{\sqrt{3}} (\bar{d}_a \sigma^{mn} \gamma^5 u_a) \bar{\partial}_m U_m^I,
\]

In order to estimate the masses of the neutral chiral bosons we can use the mass matrix \[16\]. Since the neutral chiral bosons belong to the same multiplets as their charged partners and they do not mix, their masses can be estimated as

\[
M_U = m = \sqrt{2} M_L \approx 719 \text{ GeV},
\]
\[
M_T = M = 2 M_L \approx 1017 \text{ GeV}.
\]

Degeneracies in the masses of \(U_m^I, U_m^R\) and \(T_m^I, T_m^R\) states are removed as a result of mixings of \(U_m^R\) and \(T_m^R\) states with the photon and \(Z\) boson, leading to the inequalities \(M_{\gamma R} > M_{U^I} = M_U\) and \(M_{\gamma R} > M_{T^I} = M_T\).

It is a hard task to detect neutral tensor interactions on the background electromagnetic and neutral weak processes in low energy experiments. Up to now there are no positive indications, besides, may be, the deviation in the muon anomalous magnetic moment, which can be assigned to mixings between the photon and the neutral chiral bosons. Nevertheless, as we will show in the next section the chiral bosons have a unique signature for their detection at hadron colliders due to their anomalous couplings with fermions.

### IV. RESULTS AND CONCLUSIONS

The main interest represents the production cross-sections of the lightest chiral bosons. In order to estimate them, we will calculate basic partial decay widths of the chiral bosons into leptons and quarks.

Using the Yukawa interactions \[21\] we can evaluate the lepton

\[
\Gamma_{\ell} \equiv \Gamma(\ell \tau \to \ell \nu) = \frac{g^2 M_L}{192\pi} = \frac{\Gamma(W \to \ell \nu)}{\sqrt{30}} \approx 360 \text{ MeV}
\]

and the quark

\[
\Gamma_q \equiv \Gamma(\ell \to q d) = 15 \Gamma_{\ell} \approx 5.4 \text{ GeV}
\]

widths of the charged chiral bosons \(\hat{U}^\pm\). Assuming the presence of only three fermion generations with masses lighter than \(M_L\) we can estimate the total width as \(\Gamma = 48 \Gamma_{\ell} \approx 17.2 \text{ GeV}\).

Then it is possible to evaluate the production cross-sections at the Tevatron \[26\]

\[
\sigma_{\ell}^{\text{Tev}} = \frac{4\pi^2 \Gamma_{\ell}^\tau}{3 M_L^2} \int u(x, M_L) d \left(\frac{\tau}{x}, \frac{M_L}{x}\right) \frac{dx}{x} \approx \begin{cases}
8.4 \text{ pb} & \text{RUN I} \\
11.7 \text{ pb} & \text{RUN II}
\end{cases}
\]

and the LHC

\[
\sigma_{\ell}^{\text{LHC}} = \frac{8\pi^2 \Gamma_{\ell}^\tau}{3 M_L^2} \int u(x, M_L) d \left(\frac{\tau}{x}, \frac{M_L}{x}\right) \frac{dx}{x} \approx 0.36 \text{ nb},
\]

\[
\sigma_{\ell}^{\text{LHC}} = \frac{8\pi^2 \Gamma_{\ell}^\tau}{3 M_L^2} \int u(x, M_L) d \left(\frac{\tau}{x}, \frac{M_L}{x}\right) \frac{dx}{x} \approx 0.19 \text{ nb}
\]

(using, for example, CTEQ6M parton distribution functions \[18\] and the appropriated parameter \(\tau = M_L^2/s\), where \(s\) is the square of the center-of-mass energy).

Such a big production cross-sections would immediately contradict to the present Tevatron data, unless the chiral bosons have unusual properties. It is noteworthy, that the lepton width \[24\] is suppressed in favor of the quark one \[25\]. Hence, the lightest charged chiral particle shows leptophobic property that leads to a reduced lepton production

\[
\sigma_{\ell}^{\text{Tev}} = \sigma_{\ell}^{\text{Tev}} \times B(\ell \tau \to \ell \nu) \approx \begin{cases}
0.18 \text{ pb} & \text{RUN I} \\
0.24 \text{ pb} & \text{RUN II}
\end{cases},
\]
\[
\sigma_{\ell}^{\text{LHC}} \approx 7.5 \text{ pb}, \quad \sigma_{\ell}^{\text{LHC}} \approx 4 \text{ pb}
\]

which is the main channel of charged boson detection at hadron colliders.

However, the cross sections \[25\] are still big to remain undetected at the Tevatron. This fact can be explain by another unusual and unexpected feature of the chiral bosons connected to their anomalous interactions \[10\] with fermions. Let us compare the normalized angular
distributions of the lepton from the decays of \( W^- \)

\[
\frac{dN_W}{d\Omega} = \begin{cases} 
\frac{2}{3\pi} (1 + \cos \theta)^2, & \lambda = \pm 1, \\
\frac{4}{3\pi} \sin^2 \theta, & \lambda = 0, 
\end{cases}
\] (29)

and the \( U^\pm \)

\[
\frac{dN_{U^\pm}}{d\Omega} = \begin{cases} 
\frac{2}{3\pi} \sin^2 \theta, & \lambda = \pm 1, \\
\frac{4}{3\pi} \cos^2 \theta, & \lambda = 0, 
\end{cases}
\] (30)

where \( \lambda \) is the boson helicity.

For example, the left-handed quark \( d \) (from the proton) interacting with the right-handed anti-quark \( \bar{u} \) (from the anti-proton) can produce the \( W^- \) with spin projection on the proton beam direction \(-1\). Hence, the decay leptons are distributed as \((1 + \cos \theta)^2\). While chiral particle production arises from the interaction of a quark and an anti-quark with the same helicities, that leads to zero helicity of the produced chiral boson and \( \cos^2 \theta \) lepton distribution [19].

Indeed, the cross section for \( p + \bar{p} \rightarrow U + X \rightarrow \ell + X \) process

\[
d\sigma = \frac{1}{3} \int dx_1 dx_2 u(x_1, M_L) d(x_2, M_L) d\hat{\sigma}(\hat{s}, \hat{t})
\] (31)

is expressed through the relevant differential cross section of the parton subprocess \( d + \bar{u} \rightarrow U^- \rightarrow \ell + \bar{\nu} \)

\[
E_\ell \frac{d^3\hat{\sigma}(\hat{s}, \hat{t})}{d^3p_\ell} = \frac{5g^4}{(32\pi)^2} \frac{(\hat{s} + 2\hat{t})^2 \delta(\hat{s} + \hat{t} + \hat{u})}{\hat{s} - M_L^2 + iM_L\Gamma^2}.
\] (32)

where \( \hat{s} = (p_d + p_{\bar{u}})^2, \hat{t} = (p_\ell - p_{\bar{t}})^2 \) and \( \hat{u} = (p_d - p_{\bar{u}})^2 \) are the Mandelstam variables. In the center-of-mass parton system the differential cross section shows the following system

\[
\frac{d\hat{\sigma}}{d\Omega} \propto (\hat{s} + 2\hat{t})^2 \propto \cos^2 \hat{\theta} = 1 - \frac{4\hat{p}_{\ell T}^2}{\hat{s}}.
\] (33)

Here \( \hat{\theta} \) is the angle between the lepton and the parton direction and \( \hat{p}_{\ell T}^2 \) is the square of the transverse lepton momentum.

Since the latter is invariant under longitudinal boosts along the beam direction, the distribution [33] versus \( \hat{p}_{\ell T} \) holds also in the lab frame \( p_T = \hat{p}_{\ell T} \). Changing variables in the differential cross section from \( \cos \hat{\theta} \) to \( \hat{p}_{\ell T}^2 \)

\[
\frac{d\cos \hat{\theta}}{dp_{\ell T}^2} = \frac{2}{s} \left( \sqrt{1 - \frac{4p_{\ell T}^2}{s}} \right)^{-1}
\] (34)

leads to a kinematical singularity at the endpoint \( \hat{p}_{\ell T}^2 = \hat{s}/4 \), which gives the prominent Jacobian peak in the \( W \) decay distribution.

In contrast to this, the pole in the decay distribution of the chiral bosons is cancelled out and, moreover, the distribution reaches zero at the endpoint \( p_{\ell T}^2 = \hat{s}/4 \). The chiral boson decay distribution has a broad smooth bump with the maximum below the kinematical endpoint, instead of sharp Jacobian peak (Fig. 1). In the case when the chiral boson is produced with no transverse momentum, the transverse mass of the lepton pair is related to \( p_T \) as \( M_T(\ell\nu) = 2p_T \) and the Jacobian peak is absent in \( M_T \) distribution as well.

Therefore, the decay distribution of the chiral bosons differs drastically from the distribution of the gauge bosons. Even relatively small decay width of the chiral bosons leads to a wide distribution, that obscures their identification as resonances at hadron colliders. At the present time one can speak only about eventual excess of events in the region \( 350 \text{ GeV} < M_T < 500 \text{ GeV} \), where the background from the tail of the \( W \) decays is considerably small. Indeed, such excess, although not statistically significant, has been pointed out recently by the CDF Collaboration [20] in the same region.

The form of the decay distribution for the chiral bosons resembles the bump anomalies in the inclusive jet \( E_T \) distribution, reported by the CDF Collaboration [21]. It has been suggested that these anomalies could be described by a new neutral leptophobic gauge boson with huge decay width [22]. Although this problem has been solved in the framework of the SM by changing the gluon distribution functions [23], it could be reconsidered in the light of the new form of the decay distribution as a real physical signal from decays of different chiral bosons, both charged and neutral.

Analysing the bumps in the jet transverse energy distribution in Fig. 1 of ref. [21], we can find the endpoint of the first bump at 250 GeV and guess about the second bump endpoint from the minimum around 350 GeV. If we assign the first bump to the hadron decay products of the lightest charged bosons, which exactly corresponds to the estimated mass from eq. [20], the second endpoint hints to a mass for the lightest neutral boson around 700 GeV, which is also in a qualitative agreement.
with our estimations [23]. However, taking into account the large systematic uncertainties in jet production, these conclusions may be premature.

Nevertheless, the lightest neutral chiral bosons $U_R^m$ and $U_L^m$ can fairly play the role of the hadrophobic boson, suggested in [22] in order to explain Tevatron and LEP anomalies, with nearly the same mass. Indeed, these bosons do not decay into charged leptons and cannot be seen in lepton channel. In the best case they have invisible decay widths in neutrinos, if right-handed neutrinos are not too heavy. However, in contrast to [22] these bosons are non-gauge bosons, but chiral ones. There exist also additional neutral chiral bosons $T_R^m$ and $T_L^m$, which couple only to the $u$ fermions with approximately the same Yukawa coupling constants as $U_R^m$ and $U_L^m$ bosons. They influence the low-energy hadron physics two times weaker than the $U_R^m$ and $U_L^m$ bosons due to their higher masses.

The production of the heaviest neutral chiral bosons $T_R^m$, $T_L^m$ and the detection of their leptonic decays at the hadron colliders will be crucial test of the discussed model. Unfortunately, due to their big masses they cannot be produced at the Tevatron at present. Let us calculate their lepton and quark widths in order to estimate their production cross-section at the LHC.

Using the Yukawa interactions [22] we obtain the following decay widths

$$\Gamma^0_\ell \equiv \Gamma(T^0 \to \ell^+\ell^-) = \frac{g^2 M^2_\ell}{48 \pi} \approx 2.9 \text{ GeV}$$ (35)

for the lepton and

$$\Gamma^0_q \equiv \Gamma(T^0 \to \bar{d}d) = 3 \Gamma^0_\ell \approx 8.6 \text{ GeV}$$ (36)

for the quark channels. This gives $\Gamma^0_T = 12 \Gamma^0_\ell \approx 34 \text{ GeV}$ for the total fermion decay width. Since the lepton branching ratio $B(T^0 \to \ell^+\ell^-) = 1/12$ is not negligible one expects essential signal in the dilepton channels with a cross section

$$\sigma_{\ell^+\ell^-}^{\text{HC}} = \sigma_{T^0}^{\text{HC}} \times B(T^0 \to \ell^+\ell^-)$$

$$= \frac{8\pi^2 g^2 \Gamma^0_\ell}{3 M^2_\ell} \int d\tau(x, M_T) \frac{d\tau}{x} \approx 1 \text{ pb}.$$ (37)

The most conclusive experimental evidence for the presence of anomalous interactions and new type of spin-1 non-gauge bosons will be the detection of wide bumps in leptonic distribution at high $p_T$ at Tevatron and LHC and the dilepton production with high invariant mass around 1 TeV. Considering $e^\pm$ and $\mu^\pm$ channels for the decays of the lightest charged chiral bosons, one may expect around 1 event per day at Tevatron (average luminosity $10^{34} \text{ cm}^{-2}\text{s}^{-1}$) and around 1 event every 5 seconds at LHC (designated peak $L = 10^{34} \text{ cm}^{-2}\text{s}^{-1}$). Taking into account also that in $d\bar{d}$-annihilation two different species of the neutral chiral bosons $T_R^m$ and $T_L^m$ are produced with approximately the same masses, one may expect one electron or muon pair with invariant mass around 1 TeV every 25 second at LHC peak luminosity.

In conclusion, we would like to note that our numerical results essentially depend on the magnitude of the low-energy effective tensor coupling constant $f_T$. According to our estimations [24] its value can be up to 30% larger, which effectively decreases all the calculated masses of the chiral bosons by 15%. However, the mass ratios of the chiral bosons

$$M_L : M_U : M_T : M_H = 1 : \sqrt{2} : 2 : \sqrt{5}$$ (38)

are fixed by the model. Really, the masses of the charged $U_R^m$, $T_R^m$ and neutral $U_L^m$, $T_L^m$ chiral bosons can get an additional positive contributions due to the small mixings between the gauge and the chiral bosons, which have been neglected here.

Taking into account the unusual lepton distribution [33], the main attention should be paid to events with large pseudorapidities, where the dominant lepton production occurs as a result of decays of the new chiral particles. To support these qualitative statements and to be ready to meet eventual new physics, detailed quantitative calculations should be fulfilled on the basis of the interactions [21] and [22], taking into account also all possible known effects and specifics of concrete detector.

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[25] It is interesting to note that in the same paper the author has made a conclusion about the impossibility to introduce an antisymmetric tensor in a chiral-invariant manner. It has been done in [2].
[26] Since the $\mathcal{U}$ cross section is mainly determined by a valence quark from the proton and an anti-quark from the anti-proton, we neglect here the contribution of the sea quarks.