Kinetic mixing, custodial symmetry, $Z$, $Z'$ interactions and $Z'$ production in hadron colliders.

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In this work we study the interactions of the $Z$ and $Z'$ generated by kinetic mixing in a class of theories for physics beyond the standard model motivated by the dark matter problem, containing a spontaneously broken extra $U(1)_d$ gauge factor group in a hidden scenario and a Higgs sector respecting custodial symmetry. It is shown that custodial symmetry allows us to write the $Z$ and $Z'$ couplings to standard model fermions in terms of the measured values of $\alpha, G_F, M_Z$ and $M_{Z'}$. Working at the loop level, we calculate the ratio $\rho = M_W^2/C_Z^2 M_Z^2$ used in the fit to electroweak precision data (EWPD) and use its value to estimate possible effects of kinetic mixing at the electroweak scale. For the $Z$ sector, we calculate the oblique parameters $S$ and $T$, finding that for $M_{Z'} \geq M_Z$ our results are in agreement with the values of the oblique parameters extracted from the global fit to EWPD at $1\sigma$ level. As to the $Z'$ sector, we calculate the $Z'$ contributions to charged lepton pair production at the Large Hadron Collider in the well motivated case of dark matter entering particle physics as the matter fields of the $U(1)_d$ gauge symmetry with perturbative couplings at the electroweak scale, finding that data reported by the Compact Muon Selenoid Collaboration impose a lower limit $M_{Z'} \gtrsim 5.0$ TeV.

I. INTRODUCTION

Additional $U(1)$ gauge symmetries at low energies are predicted in a variety of models for physics beyond the standard model (SM). Grand unified theories with gauge groups of rank higher than the rank of the SM gauge group, yield naturally $U(1)$ factor groups [1], [2], and there exists a classification of these possibilities which have distinctive signatures for the low energy effects of a new massive physical neutral gauge boson, usually named $Z'_r$ [3], [4], [5]. One of the most important effects at low energies of this new physics, is the generation of a kinetic mixing between the $U(1)_d$ gauge bosons. Indeed, the renormalization group flow for the coupling constants from high to low energies yields a dimension four operator $B_{\mu \nu} V_{\mu \nu}$ where $V_{\mu \nu}$ denotes the stress tensor for the $U(1)$ gauge boson $V_\mu$, even if it vanishes at some high energy scale [3], [6]. This term, modifies the expected behavior of the low energy theory based only on the charges of SM fermions in the considered ultraviolet completion.

There are two main effects of having a non-vanishing kinetic mixing at the electroweak scale. First, it produces a non-canonical form of the kinetic terms which can be diagonalized by a $GL(2, \mathbb{R})$ transformation generating small couplings of the new canonical field to the hypercharge of SM fields [5]. These new couplings triggered the interest in using kinetic mixing as an alternative to conventional mechanisms to connect the SM with one of the most challenging problems today in high energy physics, the mystery of the nature of dark matter [7], [8], [9], [10], [11], [12], [13], [14], [15], [16], [17], [18], [19], [20], [21], [22], [23], [24], [25], [26], [27], [28], [29], [30]. Second, in the presence of kinetic mixing, new mass terms are generated by the spontaneous breaking of gauge symmetries.

In the standard model, the Higgs sector has a global $SU(2)_L \otimes SU(2)_R$ symmetry which under spontaneous symmetry breaking (SSB) breaks down to its diagonal subgroup $SU(2)_V$. The $SU(2)_L$ gauge symmetry generators $T$ transform as a triplet $[31], [32], [33]$ under this residual symmetry, named custodial symmetry in the literature [33]. This global symmetry requires the corresponding gauge bosons $W_\mu$ to have a common mass and protects this property against radiative corrections. However, in the SM the $W_\mu^3$ component mixes with the gauge boson $B_\mu$ of the $U(1)_Y$ gauge symmetry to produce the massive $Z_\mu$ and the massless photon, such that the $Z_\mu$ mass is related to the $W_\mu^\pm$ mass at tree level as $M_W = M_Z \cos \theta_w$. Radiative corrections involving $U(1)_Y$ charged operators yield small corrections to this mass relation which turn out to be of the order of one percent [34].

Physics beyond the standard model may modify this picture. New Higgs fields transforming in higher $SU(2)_L$ representations break custodial symmetry and modify the custodial symmetry relations [35]. The value of the ratio $M_W^2/M_Z^2 \cos^2 \theta_w$ extracted from the global fit to electroweak precision data [31] puts stringent constraints on the possibility of custodial symmetry violating Higgs structures.

The relation $M_W = M_Z \cos \theta_w$ can also be modified by physics beyond the standard model with extra $U(1)$ gauge symmetries even if the Higgs sector respects custodial symmetry. In this case, the modification enters through the mixing of $W^3$ with more than one $U(1)$ gauge bosons. The stringent constraints on the photon mass [34] requires this mixing to preserve the unbroken nature of the electromagnetic $U(1)_{em}$ group generated by $Q = T^3 + Y/2$, in
which case, the custodial symmetry relation holds for a $\tilde{Z}_\mu$ field which is however not diagonal. The diagonalization procedure modifies the custodial symmetry relation for the physical field $Z_\mu$ and generates couplings of the physical extra gauge boson $Z_\mu'$ to SM fields. The physical outcome of this scenario depends on details such as possible charges of SM fields of the new $U(1)$ gauge symmetries and if the new fields carry SM charges. Under SSB these charges may generate new mass terms yielding a rich scenario which is however constrained by electroweak precision data.

The present work is motivated by the dark matter connection, where fields in the ultraviolet completion belong to the dark matter sector. In this case, the coupling of SM fields with the dark fields must be very small and the simplest realization is to have separate SM and dark worlds, connected only by the kinetic mixing of $U(1)_Y$ and the $U(1)$ factor subgroup of the gauge symmetry of the dark sector. This hidden dark matter scenario is an alternative realization of the Weakly Interactive Massive Particle (WIMP) idea where the small couplings arises from perturbative gauge couplings and massive mediator between dark and SM fields. The possibility of a kinetic mixing connection between dark and standard model sectors has been previously considered in the literature, but we will work out here the consequences of a physically well motivated constraint for the ultraviolet completion: custodial symmetry.

It was recently shown that if custodial symmetry is respected by the extended Higgs sector, the fact that the mass term of the $W_\mu'$ is related by this symmetry to the mass term of the $W_\mu$, can be used to write the mixing parameters entirely in terms of SM observables and the mass of the $Z_\mu'$ [39]. This procedure has the advantage that we can use precision data of the SM to constrain directly the mass of the new gauge boson instead of the mixing parameters as is conventionally done. Using the results of the global fit to EWPD for the ratio $\rho_0 = M_{Z}^{2}/M_{Z'}^{2} \overline{c}_Z \hat{\rho}$ at 1σ level [34], we did show in [36] that, in the considered framework, mixing relations straightforwardly yield the lower bound $M_{Z'} > M_{Z}$.

In this framework, the couplings of the $Z_\mu$ and $Z_\mu'$ bosons in the extended theory can also be written in terms of measured SM data and the mass of the $Z_\mu'$ boson. In the present work we do this rewriting and study the implications for the physics of the $Z_\mu$ boson and possible effects at the electroweak scale of the existence of a $Z_\mu'$ boson. As to the $Z_\mu'$ boson, effects at the electroweak scale of physics beyond the SM can in general be encoded in the oblique parameters $S$, $T$ and $U$ [37, 38, 39, 40, 41], thus we focus on the calculation of these parameters. Concerning effects of the $Z_\mu'$ at low energies, we calculate its contribution to the production of a lepton pair in hadron colliders and compare our results with experimental data obtained by the Compact Muon Selenoid (CMS) collaboration at the Large Hadron Collider (LHC) [42, 43].

Our work is organized as follows. In the next section we rewrite the couplings of the neutral gauge bosons in theories beyond the SM with an extra $U(1)$ factor group, kinetic mixing and a Higgs sector respecting custodial symmetry. Section III is devoted to the calculation of the oblique parameters and a comparison with results of the global fit to EWPD. In section IV we calculate the induced $Z_\mu'$ couplings to SM fermions, work out the $Z_\mu'$ contributions to lepton pair production at the LHC and compare with CMS results. Our conclusions are given in Section V.

II. ADDITIONAL $U(1)_d$, KINETIC MIXING AND CUSTODIAL SYMMETRY

Let us consider an extension of the SM to a group $G$ with a spontaneously broken factor Abelian gauge symmetry which we will denote as $U(1)_d$ in the following. At low energies the theory has the $G_{SM} \otimes U(1)_d$ gauge symmetry and the invariant Lagrangian including all dimension-four terms is given by

$$\mathcal{L} = \mathcal{L}_{SM}(\tilde{W}^a, \tilde{B}, \tilde{\phi}) + \mathcal{L}_{V}(V, \Phi) - \frac{\sin \chi}{2} V^{\mu \nu} B_{\mu \nu} - 2 \kappa \Phi^{*} \Phi \tilde{\phi},$$

where we use $\tilde{f}$ to distinguish the SM fields $f$ in the extended theory, $\tilde{B}^{\mu \nu}$ stands for the $U(1)_Y$ strength tensor, $V^{\mu \nu}$ denotes the strength tensor for the extra gauge boson $\tilde{V}_\mu$ of $U(1)_d$, and the complex Higgs field which spontaneously breaks this symmetry is denoted as $\Phi$. The explicit form of $\mathcal{L}_{V}(V, \Phi)$ depends on our choice of the ultraviolet completing theory, but results in this paper are independent of this choice, except for the structure of the Higgs sector, which upon spontaneous symmetry breaking must respect custodial symmetry.

The kinetic terms for the gauge bosons in the Lagrangian in Eq. (1) are

$$\mathcal{L}_{gauge}^{K} = -\frac{1}{4}(W_{\mu \nu} W^{\mu \nu} + B_{\mu \nu} B_{\mu \nu} + V^{\mu \nu} V_{\mu \nu} + 2 \sin \chi V^{\mu \nu} B_{\mu \nu}),$$

and contain a kinetic mixing term of the gauge bosons of $U(1)_Y$ and $U(1)_d$ factor groups. This term makes the kinetic Lagrangian not canonical and we must perform the following $GL(2, \mathbb{R})$ transformation on the gauge bosons to obtain properly normalized kinetic terms [5]

$$B_{\mu \nu} = \tilde{B}_{\mu \nu} - \tan \tilde{V}_{\mu \nu}, \quad V_{\mu \nu} = \sec \chi \tilde{V}_{\mu \nu}.$$

(3)
After this transformation, the kinetic terms gets the canonical form

$$L_{\text{gauge}}^K = -\frac{1}{4}(\tilde{W}^{\mu\nu}\tilde{W}_{\mu\nu} + \tilde{B}^{\mu\nu}\tilde{B}_{\mu\nu} + \tilde{V}^{\mu\nu}\tilde{V}_{\mu\nu}),$$

but it induces a coupling of the $U(1)_d$ gauge boson with the SM fields. Indeed, after this transformation the $SU(2)_L \otimes U(1)_Y \otimes U(1)_d$ covariant derivative reads

$$D'^{\mu} = \partial^{\mu} + i\tilde{g}T^a\tilde{W}_a^{\mu} + i\tilde{g}Y \frac{V}{2}\tilde{B}^{\mu} + i(g_d sec \chi Q_d - \tilde{g}_Y tan \frac{Y}{2})\tilde{V}^{\mu},$$

where $Q_d/2$ denotes the generator of $U(1)_d$, $g_d$ is the corresponding coupling constant and we use the same "tilde" notation for the SM electroweak gauge couplings, $\tilde{g}, \tilde{g}_Y$, in the extended theory.

The effect of the kinetic mixing propagates and reaches mass terms generated by the Higgs mechanism causing a mixing of the SM neutral bosons with the $\tilde{V}_\mu$ gauge boson to produce the physical $A_\mu, Z_\mu$ and a new physical boson denoted by $Z'_\mu$. Indeed, if we want to keep the $U(1)_e$ unbroken with a generator $Q = T_3 + Y/2$, we need the $\tilde{B}_\mu$, to be the SM hypercharge gauge boson which mixes with $\tilde{W}_3^\mu$ to produce the physical photon. Notice that the $\tilde{B}_\mu$ field has the same couplings to SM fields as the original $\tilde{B}_\mu$ field thus this is just a reinterpretation of the SM $U(1)_Y$ gauge boson. With the conventional weak rotation

$$
\begin{pmatrix}
\tilde{B} \\
\tilde{W}_3
\end{pmatrix} =
\begin{pmatrix}
\cos \hat{\theta}_w & -\sin \hat{\theta}_w \\
\sin \hat{\theta}_w & \cos \hat{\theta}_w
\end{pmatrix}
\begin{pmatrix}
A \\
\tilde{Z}
\end{pmatrix}
$$

we obtain

$$\tilde{g}_T \tilde{W}_3 + \tilde{g}_Y \frac{Y}{2} \tilde{B} = eQA + \frac{\tilde{g}}{\hat{c}_w}(T_3 - s^2 w Q)\tilde{Z},$$

with $e = \tilde{g}\hat{s}_w = \tilde{g}_Y\hat{c}_w$. Hereafter, we will use the shorthand notation $\hat{s}_x = \sin \hat{\theta}_x, \hat{c}_x = \cos \hat{\theta}_x$ for the mixing angles.

The Lagrangian for the Higgs sector of the $G_{\text{SM}} \otimes U(1)_d$ gauge theory reads

$$\mathcal{L}_{\text{Higgs}} = (D^{\mu}\tilde{\phi})^\dagger D^{\mu}\tilde{\phi} + (D^{\mu}\Phi)^\dagger D^{\mu}\Phi - V(\tilde{\phi}, \Phi),$$

with the following Higgs potential

$$V(\tilde{\phi}, \Phi) = \mu^2\tilde{\phi}^\dagger\tilde{\phi} + \lambda(\tilde{\phi}^\dagger\tilde{\phi})^2 + \mu^2\Phi^\dagger\Phi + \lambda_d(\Phi^\dagger\Phi)^2 + 2\kappa\Phi^\dagger\Phi\tilde{\phi}^\dagger\tilde{\phi}.$$

We are interested in the effects of kinetic mixing here, thus we will assume that the SM Higgs $\tilde{\phi}$ is a singlet under $U(1)_d$ and the new Higgs field $\Phi$ is a singlet under the SM group. In the unitary gauge, Eq. (8) yields the following gauge bosons mass terms

$$\mathcal{L}_{\text{mass}} = M_{\tilde{W}}^2 \tilde{W}_3^{\dagger}\tilde{W}_3 + \frac{1}{2} \left(M_{\tilde{Z}}^2 \tilde{Z}_3^{\dagger}\tilde{Z}_3 + 2\Delta \tilde{V}_3^{\dagger}\tilde{Z}_3 + M_{\tilde{V}}^2 \tilde{V}_3^{\dagger}\tilde{V}_3 \right),$$

with

$$M_{\tilde{W}}^2 = \frac{\tilde{g}_T^2\hat{c}_w^2}{4}, \quad M_{\tilde{Z}}^2 = \frac{M_{\tilde{V}}^2}{\hat{c}_w^2}, \quad \Delta = \frac{M_{\tilde{V}}^2}{\hat{c}_w^2} s^2 w \tan \chi, \quad M_{\tilde{V}}^2 = M_{\tilde{W}}^2 \hat{c}_w^2 \tan^2 \hat{\theta}_w \tan^2 \chi + g^2_d v_d^2 \sec^2 \chi.$$

The photon is massless, the SM field $\tilde{Z}$ has the expected mass value from custodial symmetry but a mixing with the $\tilde{V}$ field has been generated by the kinetic mixing. The neutral gauge boson part of this Lagrangian can be diagonalized by the following rotation

$$
\begin{pmatrix}
\tilde{Z} \\
\tilde{V}
\end{pmatrix} =
\begin{pmatrix}
\cos \hat{\theta}_\zeta & -\sin \hat{\theta}_\zeta \\
\sin \hat{\theta}_\zeta & \cos \hat{\theta}_\zeta
\end{pmatrix}
\begin{pmatrix}
\tilde{Z}' \\
\tilde{V}'
\end{pmatrix}
$$

After these transformations, the original gauge fields are related to the diagonal fields by the following matrix

$$
\begin{pmatrix}
\tilde{B} \\
\tilde{W}_3 \\
\tilde{V}
\end{pmatrix} =
\begin{pmatrix}
\hat{c}_w, -\hat{s}_w c_\zeta - \tan \chi s_\zeta, -\hat{c}_w \hat{s}_w s_\zeta - \tan \chi c_\zeta \\
\hat{s}_w & \hat{c}_w c_\zeta & -\hat{c}_w s_\zeta \\
0 & \sec \hat{\chi} s_\zeta & \sec \hat{\chi} c_\zeta
\end{pmatrix}
\begin{pmatrix}
A \\
\tilde{Z} \\
\tilde{Z}'
\end{pmatrix}.
$$
and in terms of the physical fields the covariant derivative reads

\[ D_\mu = \partial_\mu + i \frac{\tilde{g}}{\sqrt{2}} (T^+ \tilde{W}_\mu^+ + T^- \tilde{W}_\mu^-) + i e Q A_\mu \]

\[ + i \left[ \frac{\tilde{g} c_\chi}{c_w} \left( (T_3 - s_w^2 Q) - \tilde{s}_w \tan \theta_\chi \tan \frac{\chi}{2} \right) + g_d s_\chi \sec \chi \frac{Q_d}{2} \right] Z_\mu \]

\[ - i \left[ \frac{\tilde{g} s_\chi}{c_w} \left( (T_3 - s_w^2 Q) + \tilde{s}_w \tan \theta_\chi \tan \frac{\chi}{2} \right) - g_d c_\chi \sec \chi \frac{Q_d}{2} \right] Z'_\mu. \] (14)

The effects of kinetic mixing are conventionally analyzed comparing experimental data with predictions from the neutral currents arising from this covariant derivative. This comparison yields usually bounds of the possible values of the kinetic mixing parameter \( \chi \), the mixing angle \( \theta_\chi \) or the mixing angle in the Higgs sector due to the \( \kappa \) term in Eq. (9), not shown here. In a recent work \[36\] we pointed out that the tree level custodial symmetry protected relation \( M_Z^2 = M_W^2 \cos^2 \theta_w \), allows to write the matrix elements in Eq. (13) in terms of the measured weak angle, \( M_W \), \( M_Z \) and the unknown mass of the diagonal field \( Z' \). Aiming to use results of the fit to electroweak precision data which requires a loop level analysis, we argue that loop contributions are dominated by the effects of standard model particles, and the main effect of kinetic mixing in the mass Lagrangian can be taken at the tree level.

Effects of new physics are considered in the global fit to the EWPD through the parameter \[34\]

\[ \rho_0 \equiv \frac{M_W^2}{c_Z^2 M_Z^2 \tilde{\rho}}, \] (15)

where \( \tilde{c}_Z \equiv \cos \theta_w(M_Z) \) is the Weinberg angle measured at the \( Z \) pole and \( M_W^2 \), \( M_Z^2 \) denote the masses of the physical \( Z \) and \( W^\pm \) respectively. The quantity \( \tilde{\rho} \) in Eq. (15) accounts for radiative corrections in the SM such that if there are no new physics contributions then \( \rho_0 = 1 \). Deviations from this value are necessarily due to physics beyond the SM.

Radiative corrections to the ratio \( M_W^2/c_Z^2 M_Z^2 \) due to SM particles are dominated by the top quark and the next to leading order effects are given by Higgs boson loops. In the \( \overline{\text{MTS}} \) scheme, including all boson contributions yields \[34\]

\[ \tilde{\rho} = 1.01019 \pm 0.00009. \] (16)

The global fit to electroweak precision data yields \[34\]

\[ \rho_0 = 1.00038 \pm 0.00020. \] (17)

Notice that \( \rho_0 > 1 \) at 1.9 \( \sigma \) (94% confidence level) \[34\]. Although not conclusive, this result points to possible new physics contributions to the value of \( \rho_0 \) which we will assume to be dominated by the kinetic mixing effects. In the following we will explore the consequences of having \( \rho_0 \neq 1 \), aiming to constrain the possible values of the \( Z' \) mass from available data.

Comparison with electroweak precision data requires to calculate \( \rho_0 \) in the extended theory. Since in the hidden \( U(1)_d \) extension of the standard model \( M_W \) is the physical mass of the \( W^\pm \) we have

\[ \rho_0 = \frac{M_W^2}{c_Z^2 M_Z^2 \tilde{\rho}}. \] (18)

Using custodial symmetry relations at the loop level it has been shown that the \( \tilde{Z} - \tilde{V} \) mixing angle in the \( \overline{\text{MTS}} \) scheme is given by \[36\]

\[ \tilde{s}_\chi^2 \equiv \frac{\sigma_0 (\rho_0 - 1)(\rho_0 c_Z^2 - \tilde{s}_Z^2)}{(\rho_0 - \tilde{s}_Z^2)(\rho_0 - \sigma_0)}, \] (19)

where \( \tilde{s}_\chi = s_\chi(M_Z) \), is the sine of \( \tilde{Z} - \tilde{V} \) mixing angle at the scale \( \mu = M_Z \) and

\[ \sigma_0 \equiv \frac{M_W^2}{c_Z^2 M_Z^2 \tilde{\rho}}. \] (20)

Similarly, the kinetic mixing angle at the same scale, \( \tilde{\chi} = \chi(M_Z) \), can be written in terms of the same physical quantities as

\[ \tan^2 \tilde{\chi} = \frac{(\rho_0 - 1)(\rho_0 c_Z^2 - \tilde{s}_Z^2)}{\rho_0 \sigma_0 \tilde{s}_Z^2 c_Z} \frac{1}{\rho_0 (1 - \sigma_0 c_Z^2) - \tilde{s}_Z^2}. \] (21)
Using these relations, the covariant derivative in Eq. (14) can be rewritten in terms of the physical values of the effective couplings depending on \( \alpha, G_F, M_W, M_Z \) and \( M_Z \) as

\[
D_\mu = \partial_\mu + ie\frac{\sqrt{\rho_0}}{\sqrt{2}s_Z}(T^+W^+_\mu + T^-W^-_\mu) + ieQA_\mu
\]

\[
\quad + i \left[ \frac{e}{s_Zc_Z} \sqrt{\frac{\rho_0 - s_Z^2 - \rho_0\sigma_0c_Z^2}{c_Z^2\rho_0(\rho_0 - \sigma_0)}} (T_3 - (1 - \rho_0c_Z^2)Q) + g_d\hat{s}_c\sec\hat{\chi} \frac{Q_d}{2} \right] Z_\mu
\]

\[
\quad - i \left[ \frac{e}{s_Zc_Z} \sqrt{\frac{(\rho_0 - 1)(\rho_0c_Z^2 - s_Z^2)}{c_Z^2\sigma_0(\rho_0 - \sigma_0)}} (T_3 - (1 - \sigma_0c_Z^2)Q) - g_d\hat{c}_c \sec\hat{\chi} \frac{Q_d}{2} \right] Z'_\mu,
\]

where

\[
\hat{c}_c^2 = \frac{\rho_0(\rho_0 - s_Z^2 - \rho_0\sigma_0c_Z^2)}{(\rho_0 - \sigma_0)},
\]

\[
\hat{s}_c \sec\hat{\chi} = \frac{1}{\rho_0s_Zc_Z^2} \left[ \frac{(\rho_0 - 1)(\rho_0c_Z^2 - s_Z^2)}{(\rho_0 - 1)(\rho_0c_Z^2 - s_Z^2)} + \frac{\rho_0\sigma_0c_Z^2}{\rho_0(\rho_0 - \sigma_0)} \right]^{\frac{1}{2}},
\]

\[
\hat{c}_c \sec\hat{\chi} = \frac{1}{s_Zc_Z} \left[ \frac{\rho_0 - s_Z^2 - \rho_0\sigma_0c_Z^2}{\rho_0 - \sigma_0} \left( \frac{\rho_0 - 1}{\rho_0(\rho_0 - \sigma_0)} \right)^2 + \rho_0s_Z^2 \right]^{\frac{1}{2}},
\]

with \( \hat{c}_c = c_c(M_Z) \).

III. EFFECTIVE \( Z\bar{f}f \) INTERACTIONS AND OBLIQUE PARAMETERS

The effective Lagrangian formulation for physics beyond the SM encode corrections to the SM Lagrangian in the oblique parameters \( S, T \) and \( U \). For the \( Z\bar{f}f \) interaction the effective Lagrangian reads \( 5, 44, 45, 46, 47 \)

\[
\mathcal{L}_{Z\bar{f}f}^{eff} = \frac{e}{2s_Wc_W} \left( 1 + \frac{\alpha T}{2} \right) \sum_f \bar{f}\gamma^\mu \left( T^3_fL - 2s_W^2c_W^2\gamma^5 \right) fZ_\mu,
\]

where \( s_W = \sin \theta_W, c_W = \cos \theta_W \) with \( \theta_W \) standing for the value of the Weinberg angle at the considered low energy scale and

\[
s^2 = s^2 + \frac{1}{c_W - s_W} \left( \frac{\alpha S}{4} - s^2c_W\alpha T \right).
\]

If the SM fermions do not carry \( U(1)_d \) charge, the covariant derivative in Eq. (22) yields the following \( Z\bar{f}f \) Lagrangian at the scale \( \mu = M_Z \)

\[
\mathcal{L}_{Z\bar{f}f} = \frac{e}{2s_Zc_Z} R \sum_f \bar{f}\gamma^\mu \left[ T^3_fL - 2(1 - \rho_0c_Z^2)Q - T^3_fL\gamma^5 \right] fZ_\mu,
\]

with

\[
R = \sqrt{\frac{\rho_0 - s_Z^2 - \rho_0\sigma_0c_Z^2}{c_Z^2\rho_0(\rho_0 - \sigma_0)}}.
\]

A comparison of Lagrangians in Eqs. (26,28) yields

\[
\alpha S = 4c_Z^2 \left[ (1 - \rho_0)(\hat{c}_c^2 - s_Z^2) + 2s_Z^2(R - 1) \right],
\]

\[
\alpha T = 2(R - 1).
\]

Notice that although \( S \) and \( T \) are functions of the \( Z' \) mass, there is always a linear relation among these parameters

\[
T = \frac{1}{4s_Z^2c_Z^2} S + (\rho_0 - 1) \frac{c_Z^2 - s_Z^2}{\alpha s_Z^2}.
\]
valid for all values of $M_{Z'}$. The values extracted from the fit to EWPD for the oblique parameters are

$$S = -0.01 \pm 0.10, \quad T = 0.03 \pm 0.12, \quad U = 0.02 \pm 0.11.$$  \hspace{1cm} (33)

In Fig. (1) we plot the predicted values of $S$ and $T$ as functions of the $Z'$ mass, for the range of values for $\rho_0$ extracted from the fit to EWPD in Eq. (33). We also show in these plots the 1σ regions for $S$ and $T$ obtained in the fit. We notice first that $S$ and $T$ reach a saturation value for $M_{Z'} \approx 250$ GeV and are not sensitive to the value of $M_{Z'}$ beyond this point. The predicted values of $S$ and $T$ are consistent with results from the fit to EWPD for $M_{Z'} > M_Z$. The linear relation between $T$ and $S$ is shown in Fig. (2) together with the 1σ bands for the oblique parameters. The intersection of the three bands yields the values of $S$ and $T$ for which the predictions of the present formalism agrees with the fit to EWPD at 1σ level.

![Figure 1: Oblique parameters $S$ and $T$ as functions of $M_{Z'}$. The solid blue lines corresponds to the predictions using the central value of $\rho_0$ and the shadow band to the 1σ region for $\rho_0$ in Eq. (17). The red bands correspond to the 1σ region for $S$ and $T$ in Eq.(33). Solid lines correspond to the central values.](image)

**IV. Z' CONTRIBUTIONS TO CHARGED LEPTON PAIR PRODUCTION AT HADRON COLLIDERS**

**A. General formalism**

Upper bounds have been obtained for the $Z'$ contributions to charged lepton pair production at the Tevatron [48] and the LHC [42], [43]. The production cross section for a fermion pair in hadron colliders in the $Z'$ pole region can be written in general as [49], [50]

$$\sigma_{\bar{f}f} = \int_{(M_{Z'}, -\Delta)^2}^{(M_{Z'}, +\Delta)^2} \frac{d\sigma}{dM^2} (pp \rightarrow Z'X \rightarrow \bar{f}fX) dM^2. \hspace{1cm} (34)$$

In the narrow width approximation for the $Z'$ it can be written as [49]

$$\sigma_{\bar{f}f} \approx \left( \frac{1}{3} \sum_q \frac{dL_{\bar{q}q}}{dM_{Z'}^2} \delta(\bar{q}q \rightarrow Z') \right) BR(Z' \rightarrow \bar{f}f), \hspace{1cm} (35)$$

where $\frac{dL_{\bar{q}q}}{dM_{Z'}^2}$ stands for the parton luminosities and the branching ratio for the $\bar{f}f$ channel is given by

$$BR(Z' \rightarrow \bar{f}f) = \frac{\Gamma(Z' \rightarrow \bar{f}f)}{\Gamma_{Z'}} \hspace{1cm} (36)$$
where $\Gamma_{Z'}$ denotes the total $Z'$ width.

In general, the Lagrangian for the interaction of the $Z'$ with standard model fermions can be written as [50]

\[ L_{Z'\bar{f}f} = g'Z'\mu \sum_f \bar{f}\gamma^\mu \left[ g_V^f - g_A^f \gamma^5 \right] f, \]  



such that the peak cross section reads

\[ \hat{\sigma}(\bar{q}q \rightarrow Z') = \frac{\pi g'^2}{12} \left( (g_V^q)^2 + (g_A^q)^2 \right), \]  



and the width in the $\bar{f}f$ channel is calculated as

\[ \Gamma(Z' \rightarrow \bar{f}f) = N_c \frac{g'^2 M_{Z'}}{4\pi} \left[ (g_V^f)^2 + (g_A^f)^2 + O(M_Z'^2) \right], \]  



where $N_c = 3$ for quarks and $N_c = 1$ for leptons. Neglecting the SM fermion masses, if the couplings are generation independent, the total width to fermions is given by

\[ \Gamma_{Z'} = \frac{g'^2 M_{Z'}}{4\pi} \left( 9 ((g_V^q)^2 + (g_A^q)^2 + (g_A^d)^2) + 3 ((g_V^q)^2 + (g_A^q)^2 + (g_A^\nu)^2 + (g_A^e)^2) \right). \]  

Taking into account that the top quark channel opens at $2m_t = 350$ GeV, reduces this width 18% below the $tt$ threshold, but already for $M_{Z'} = 500$ GeV this phase space correction is of the order of 2%, thus we can safely neglect fermion masses.

The cross section for the $Z'$ contributions to $l^+l^-$ production in hadron colliders can be written as [49], [50]

\[ \sigma_{l^+l^-} = \frac{\pi}{488} \left[ c_u w_u(s, M_{Z'}^2) + c_d w_d(s, M_{Z'}^2) \right], \]  



where the functions $w_{u,d}(s, M_{Z'}^2)$ depend only on the invariant Mandelstam variable $s$ of the collision and on the $Z'$ mass and are the same for any model containing neutral gauge bosons with generation-independent couplings to quarks. Explicit expressions for these functions in terms of the parton distribution functions of the colliding hadrons can be found in [49]. The coefficients $c_{u,d}$ depend on the $Z'$ couplings to fermions as

\[ c_u = \frac{g'^2}{2} \left( (g_V^q)^2 + (g_A^q)^2 \right) BR(Z' \rightarrow l^+l^-), \]  

\[ c_d = \frac{g'^2}{2} \left( (g_V^d)^2 + (g_A^d)^2 \right) BR(Z' \rightarrow l^+l^-). \]
Using the form given in [49] for the functions $w_{u,d}(s, M^2_{Z'})$, experimental data on $l^+l^-$ production in hadron colliders can be used to impose upper bounds on the intermediate $Z'$ contributions which translates into exclusion curves in the $c_u-c_d$ plane for given values of $M_{Z'}$. In our formalism the $Z'$ couplings to standard model fermions are fixed by known data and the $Z'$ mass, thus we are able to calculate the predicted values of $c_u,c_d$ as functions of $M_{Z'}$ and to compare these results with the exclusion curves extracted from the data.

B. $Z'$ contributions to $l^+l^-$ production for theories with kinetic mixing and custodial symmetry

In the case that SM fermions do not carry the $U(1)_d$ charges, the covariant derivative in Eq. (22) yields the following $Z'\bar{f}f$ interactions

$$\mathcal{L}_{Z'\bar{f}f} = g_{Z'} \sum_f \bar{f} \gamma^\mu \left[ T^3_{fL} - 2(1 - \sigma_0 c^2_Z)Q - T^3_{fL} \gamma^5 \right] f Z'_\mu, \quad (44)$$

where

$$g_{Z'} = \frac{e}{2s_Z c_Z} \sqrt{\frac{(\rho_0 - 1)(\rho_0 c^2_Z - s^2_Z)}{c^2_Z \sigma_0(\rho_0 - \sigma_0)}}. \quad (45)$$

Notice that this coupling scales as $\sqrt{(\rho_0 - 1)/\sigma_0}$, thus being small for $Z'$ masses close to the electroweak scale, but the $\sqrt{1/\sigma_0}$ factor enhance it for $M_{Z'} \gg M_W$ and eventually becomes large for large values of the $Z'$ mass. In Fig. 3 we show the associated fine structure constant $g_{Z'}/4\pi$ for the window of values at the $1\sigma$ level for $\rho_0$ in Eq. (17), which shows that for $M_{Z'} \sim 30$ TeV we enter in a non-perturbative regime. The couplings of the $Z'$ to fermions are generation independent but non-universal, depending of the $T^3_L$ and $Q$ quantum numbers of the specific fermion.

![Fine structure constant for the coupling $g_{Z'}$ induced by kinetic mixing as a function of the $Z'$ mass. The solid line corresponds to the predictions using the central value of $\rho_0$ and the shadow band to the $1\sigma$ region for $\rho_0$ in Eq. (17).](image)

Comparing the interacting Lagrangian in Eq. (44) with the general Lagrangian in Eq. (37), we obtain $g' = g_{Z'}$, which according to Eq. (45) depends only on known data and $M_{Z'}$. We also identify the following vector and axial factors

$$g^V_f = T^3_f - 2(1 - \sigma_0 c^2_Z)Q_f, \quad g^A_f = T^3_f. \quad (46)$$

These couplings yield the following total $Z'$ decay width into SM fermions

$$\Gamma_{Z'} = \frac{\alpha M_{Z'}}{4 s^2_Z c^2_Z} \left[ \frac{(\rho_0 - 1)(\rho_0 c^2_Z - s^2_Z)}{c^2_Z \sigma_0(\rho_0 - \sigma_0)} \right] \left[ 1 - 2(1 - \sigma_0 c^2_Z) + \frac{8}{3}(1 - \sigma_0 c^2_Z)^2 \right]. \quad (47)$$
The general formalism for the calculation of the $Z'$ production in hadron colliders uses the narrow width approximation \[49, 50\] and we must ensure that we are in this regime for the energies at which data is obtained. In a first approximation we will assume that the total width of the $Z'$ is given by its decays to SM fermions and will study below possible modifications to this picture. Under this assumption, $\Gamma_f^{Z'}$ in Eq.\((47)\) is the total width. In Fig. 4 we plot the ratio $\Gamma_f^{Z'}/M_{Z'}$ as a function of $M_{Z'}$ and we can see that the narrow width approximation is well satisfied up to masses of the order of 10 TeV. At the $M_{Z'} = 6$ TeV this ratio is at the 3% level reaching values of the order of 10% for $M_{Z'} = 10.6$ TeV, thus we can use safely the calculations based on the narrow width approximation up to this energy. The branching ratio for the $l^+l^-$ channel is obtained as

\[
BR(Z' \rightarrow l^+l^-) = \frac{1}{8} \frac{1 - 4(1 - \sigma_0 c_{Z}^2) + 8(1 - \sigma_0 c_{Z}^2)^2}{3 - 6(1 - \sigma_0 c_{Z}^2) + 8(1 - \sigma_0 c_{Z}^2)^2}.
\]  

(48)

Finally, a calculation of the coefficients $c_{u,d}$ in Eqs.\((42,43)\) with $g' = g_{Z'}$ and the vector and axial factors in Eq. \((46)\) yields

\[
c_u = \frac{\pi \alpha}{36 s_{Z}^2 c_{Z}^2} \left( \rho_0 - 1 \right) \left( \rho_0 c_{Z}^2 - s_{Z}^2 \right) \frac{c_{Z}^2 \sigma_0 (\rho_0 - \sigma_0)}{c_{Z}^2 \sigma_0 (\rho_0 - \sigma_0)} \left[ 9 - 24(1 - \sigma_0 c_{Z}^2)^2 + 32(1 - \sigma_0 c_{Z}^2)^2 \right] BR(Z' \rightarrow l^+l^-),
\]

(49)

\[
c_d = \frac{\pi \alpha}{36 s_{Z}^2 c_{Z}^2} \left( \rho_0 - 1 \right) \left( \rho_0 c_{Z}^2 - s_{Z}^2 \right) \frac{c_{Z}^2 \sigma_0 (\rho_0 - \sigma_0)}{c_{Z}^2 \sigma_0 (\rho_0 - \sigma_0)} \left[ 9 - 12(1 - \sigma_0 c_{Z}^2)^2 + 8(1 - \sigma_0 c_{Z}^2)^2 \right] BR(Z' \rightarrow l^+l^-).
\]

(50)

Notice that these factors satisfy the linear relation

\[
c_u = \frac{9}{9 - 12(1 - \sigma_0 c_{Z}^2)^2 + 8(1 - \sigma_0 c_{Z}^2)^2} c_d
\]

(51)

independently of the value of $BR(Z' \rightarrow l^+l^-)$. In general the proportionality coefficient depends on $M_{Z'}^2$, but for large $Z'$ mass, it reaches an asymptotic constant value to yield

\[
c_u = \frac{17}{9} c_d.
\]

(52)
In this limit, the branching ratio in Eq. (48) reaches a saturation value $BR(l^+l^-) = 1/8$ and the values of the $c_{u,d}$ factors grow like $M_Z^2$:

$$c_u \approx \frac{17}{8} \frac{\pi\alpha}{36s_Z^2c_Z^2} \left( \frac{\rho_0 c_Z^2 - s_Z^2}{\rho_0} \right) \left( \frac{\rho_0 - 1}{c_Z^2} \right) = 1.67 \times 10^{-6} \frac{M_{Z'}}{M_W},$$

$$c_d \approx \frac{5}{8} \frac{\pi\alpha}{36s_Z^2c_Z^2} \left( \frac{\rho_0 c_Z^2 - s_Z^2}{\rho_0} \right) \left( \frac{\rho_0 - 1}{c_Z^2} \right) = 0.49 \times 10^{-6} \frac{M_{Z'}}{M_W}.$$

In practice, large $Z'$ mass means $\sigma_0 \ll 1$, which is satisfied for $M_{Z'} \gtrsim 500$ GeV.

Experimental data on the upper bounds for $Z'$ production at the LHC has been translated into exclusion curves in the $c_d - c_u$ plane for given values of $M_{Z'}$. In Fig. 5 we show the exclusion curves obtained in [43] and our results. The dots in this plot are the values of $(c_d(M_{Z'}), c_u(M_{Z'}))$ from Eqs. (49,50) for the values of $M_{Z'}$ corresponding to the exclusion curves and are computed taking the central value of $\rho_0$ in Eq. (17). The dot for a given mass is marked in the same color as the corresponding exclusion curve. Uncertainties in this plot correspond to the 1 $\sigma$ region for $\rho_0$ in Eq. (17). For $M_{Z'} < 5200$ GeV the predicted values for $c_u, c_d$ including uncertainties are above the exclusion curves thus, at 1$\sigma$ level, the predicted values of the parameters are inconsistent with data.

![FIG. 5: Exclusion curves for the $c_u, c_d$ couplings extracted from Ref. [43] and the corresponding predictions in theories for physics BSM containing an extra spontaneously broken $U(1)_d$ and respecting custodial symmetry. Each dot corresponds to the predicted value for a given $Z'$ boson mass from Eqs. (49,50) and is marked in the same color as the corresponding exclusion curve. Uncertainties correspond to the 1$\sigma$ region for $\rho_0$ in Eq. (17).](image)

A more precise calculation requires to take into account the uncertainties in $c_u, c_d$ induced by changes in $BR(Z' \rightarrow l^+l^-)$ due to contributions of other decay channels to the total decay width. First, although the mixing with the SM fields generates couplings of the $Z'$ to $ZZ, W^+W^-$ and $ZH$, these couplings are proportional to single mixing factors $s_Z$ and the corresponding decay widths are proportional to $\rho_0 - 1$, thus small compared to the decay widths to fermions which are enhanced by the $1/\sigma_0$ factor and turn out to be proportional to $(\rho_0 - 1)M_{Z'}^2/M_W^2$. Second, the decay to non-SM particles in the ultraviolet completing theory may be more important since, as we can see from Eqs. (22,25), the corresponding coupling has the same enhancement factor as the coupling to fermions and we must have at least a rough estimate of the decay width to these non-SM particles. In this concern, the most important physical case today is the possibility that dark matter enter particle physics coupled to SM fields via kinetic mixing, in which case, it is natural to expect dark matter particles with masses of the order of the electroweak scale.
In order to have an estimate of these effects we consider the possibility of extra fermion field $\psi$ as (dark) matter field with $U(1)_d$ charge $Q_{\psi}^d = 2$, such that its coupling to the $V^\mu$ field in Eq. (1) is $g_v$. In this case, the $Z'\bar{\psi}\psi$ interaction arising from the covariant derivative in Eq. (22) is

$$\mathcal{L}_{Z'\bar{\psi}\psi} = -g_v \hat{c}_d \sec \hat{\chi} \bar{\psi} \gamma^\mu \psi Z'_\mu,$$

which yields the following decay width

$$\Gamma(Z' \rightarrow \bar{\psi}\psi) = \frac{\alpha_v M_{Z'}}{3 \hat{s}_Z^2 \hat{c}_Z^2} \left[ \frac{\rho_0 - \hat{s}_Z^2 - \rho_0 \hat{c}_Z^2}{\rho_0 - \sigma_0} \left( \frac{\rho_0 - 1}{\sigma_0} \left( \frac{\rho_0(1 - \sigma_0) \hat{c}_Z^2 - \hat{s}_Z^2}{\rho_0 \hat{c}_Z^2} + \rho_0 \hat{s}_Z^2 \right) \right) \right]
\times \left( 1 + \frac{2 M_{\psi}^2}{M_{Z'}^2} \right) \sqrt{1 - \frac{4 M_{\psi}^2}{M_{Z'}^2}},$$

where $\alpha_v = g_v^2/4\pi$ is the $U(1)_d$ fine structure constant. This width depends on the unknown $\alpha_v$, $M_\psi$ and $M_{Z'}$ and it is not possible to fix it with certainty. However, a reasonable estimate of its size can be obtained considering that the $U(1)_d$ interaction is perturbative and, at least for the dark matter context we can consider $M_\psi$ of the order of the electroweak scale. For larger masses, phase space reduces for a given value of the $Z'$ mass and we expect a smaller decay width. In Fig. 6 we show the decay width for $Z' \rightarrow \bar{\psi}\psi$ taking $\alpha_v = \alpha$ and $M = 100$ GeV, as a function of $M_{Z'}$, together with the result of the decay width to fermions in Eq.(47). It is clear from this plot that around $M_{Z'} = 5.2$ TeV this contribution is of the same order as the decay width to SM fermions, thus we expect the values of $(c_d, c_u)$ to decrease roughly by a factor of 1/2. Taking this contribution into account we find that the narrow width approximation is valid up to energies of 10 TeV, where the width to mass ratio reaches the 10% level. Considering the 1/2 factor correction to our previous calculation and the uncertainties in the value of $\rho_0$ in Eq.[17], from the exclusion curves given in [43] we refine the lower limit to $M_{Z'} \gtrsim 5.0$ TeV.

![FIG. 6: Estimate of the decay width of the $Z'$ to extra fermions with a mass of the order of the electroweak scale and a perturbative coupling.](image)

V. CONCLUSIONS

In this work we study the implications of kinetic mixing for the $Z$ and $Z'$ interactions in a class of models for physics beyond the SM motivated by the dark matter problem, whose gauge group contains a spontaneously broken $U(1)_d$ factor group and whose Higgs sector respects custodial symmetry. For these theories, the kinetic mixing generates a coupling of the extra neutral physical gauge boson $Z'$ to SM fermions. In a hidden dark matter scenario, custodial
symmetry allows to write the kinetic mixing parameters in terms of the measured values of $M_Z$, the electromagnetic constant $\alpha$, the Fermi constant $G_F$ and the mass of the $Z'$ [36], in such a way that the $Z$ and $Z'$ couplings to SM fermions can be written in terms of these parameters.

On the other hand, the fit to electroweak precision data requires the calculation of observables at the loop level and encode possible effects of physics beyond the standard model in the deviation of the $\rho$ from the unit value. The value of $\rho$ is fixed to $\rho_0 = 1.00038 \pm 0.00020$ from the global fit to electroweak precision data [34] thus $\rho_0 > 1$ at 1.9 $\sigma$ level (94% confidence level) and although not conclusive, present data points to the existence of new physics contributions to the $\rho_0$ parameter.

Custodial symmetry protects the equal mass relation for $W$, triplet from radiative corrections. This and the unbroken nature of the electromagnetic $U(1)_{\text{em}}$ group generated by $Q = T_3 + Y/2$ allows to relate the $Z$ add $W^\pm$ masses at the loop level in the standard model as $M_W^2 = M_Z^2/2 \hat{\rho}$, where $\hat{\rho}$ are small radiative corrections arising from custodial symmetry violating Yukawa interactions which enter the mass relations due to the mixing (characterized by the weak mixing angle) of the $W^2_\mu$ with the $U(1)_Y$ gauge boson $B_\mu$ to produce the standard model $Z_\mu$ and the photon field $A_\mu$.

In an extension of the standard model with a spontaneously broken $U(1)_d$ factor dark gauge group and kinetic mixing, if the Higgs sector respects custodial symmetry, then small couplings between the SM and dark sectors are generated, new mixing terms appear and mixing pattern in the neutral gauge sector is more involved. However, in the hidden dark matter scenario, custodial symmetry still relates the $W^3_\mu$ mass term to the SM $Z_\mu$ (denoted $\tilde{Z}_\mu$ here) mass term, although $\tilde{Z}_\mu$ is not the physical field. We argue that at the $\mu = M_Z$ scale, radiative corrections are dominated by SM particles and use this relation to calculate at the loop level the $Z$ and $Z'$ couplings to SM fermions. This and the analogous ratio $\sigma_0 = M_W^2/\hat{\rho}^2$ are small radiative corrections arising from custodial symmetry violating Yukawa interactions which enter the mass relations due to the mixing (characterized by the weak mixing angle) of the $W^2_\mu$ with the $U(1)_Y$ gauge boson $B_\mu$ to produce the standard model $Z_\mu$ and the photon field $A_\mu$.

For the $Z$ boson we calculate the oblique parameters $S$ and $T$, which encode the low energy effects of physics beyond the SM for a wide class of theories, including the ones considered here. We find that for $M_{Z'} > 200$ GeV the oblique parameters are not sensitive to the $Z'$ mass and a comparison with the values of $S$ and $T$ extracted from the global fit to electroweak precision data [34] at 1$\sigma$ level, yields results in agreement with the fit for $M_{Z'} > M_Z$. This lower bound is consistent with results found in [36] based on the physical range of values for the mixing angles.

As to the $Z'$ physics we study the intermediate $Z'$ contributions to the production of a charged lepton pair at the LHC. The corresponding cross section can written in terms of two parameters, $c_u$ and $c_d$, carrying all the information of the $Z'$ couplings to SM fermions. There are two factors in $c_u$ and $c_d$, the first of which contains only the couplings to SM fermions. The second factor is the branching ratio $BR(Z' \to l^+l^-)$ which contains information on SM and beyond the SM physics. In a first approximation, we calculate the $c_u$ and $c_d$ parameters in our formalism considering only the $Z'$ couplings to SM fermions generated by the kinetic mixing. In this case, the corresponding cross section depends on the unknown $Z'$ mass and known SM parameters. We compare the results of our calculation using the range of values of $\rho_0$ extracted from the global fit to the EWPD at the 1$\sigma$ level with the exclusion curves in the $c_u - c_d$ plane for $Z'$ masses in the range $3.8 - 7.0$ TeV obtained by the CMS Collaboration [12, 13]. This comparison shows that consistency of our calculation with the CMS data requires $M_{Z'} \geq 5.2$ TeV. This result is modified when we consider the coupling of the $Z'$ to kinematically allowed non-SM particles which enter $c_u$ and $c_d$ through the branching ratio $BR(Z' \to l^+l^-)$. We estimate these contributions in the well motivated case of dark matter entering particle physics as the matter fields of the $U(1)_d$ gauge symmetry with perturbative couplings at the electroweak scale, finding that our results for $c_u$ and $c_d$ get modified roughly by a factor of 1/2. Taking into account this correction we obtain that consistency with the CMS data requires $M_{Z'} \geq 5.0$ TeV. These lower limits are obtained considering the 1$\sigma$ region (68% confidence level) for the values of $\rho_0$ extracted from the fit to EWPD. Although present data points to the existence of new physics contributions to the $\rho_0$ parameter, definitive conclusions must await for more precise electroweak data which will eventually allow to lower the uncertainty in this parameter.
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