The Stueckelberg Extension and Milli Weak and Milli Charged Dark Matter

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Abstract. A overview is given of the recent developments in the \( U(1)_X \) Stueckelberg extensions of the Standard Model and of MSSM where all the Standard Model particles are neutral under the \( U(1)_X \), but an axion which is absorbed is charged under both \( U(1)_X \) and \( U(1)_Y \) and acts as the connector field coupling the Standard Model sector with the Stueckelberg sector. Coupled with the usual Higgs mechanism that breaks the \( SU(2)_L \times U(1)_Y \) gauge symmetry, this scenario produces mixings in the neutral gauge boson sector generating an extra \( Z' \) boson. The couplings of the extra \( Z' \) to the Standard Model particles are milli weak but its couplings to the hidden sector matter, defined as matter that couples only to the gauge field of \( U(1)_X \), can be of normal electro-weak strength. It is shown that such extensions, aside from the possibility of leading to a sharp \( Z' \) resonance, lead to two new types of dark matter: milli weak (or extra weak) and milli charged. An analysis of the relic density shows that the WMAP-3 constraints can be satisfied for either of these scenarios. The types of models discussed could arise as possible field point limit of certain Type IIB orientifold string models.

Keywords: \( U(1) \) extension, Stueckelberg, milli weak, milli-charged, dark matter

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INTRODUCTION

Through a Stueckelberg mechanism an Abelian gauge boson develops mass without the benefit of a Higgs mechanism (For the early history of the Stueckelberg mechanism see, [1–4]). Thus consider the Lagrangian

\[
\mathcal{L}_0 = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{1}{2} (m A_\mu + \partial_\mu \sigma) (m A^\mu + \partial^\mu \sigma),
\]

which is gauge invariant under the transformations \( \delta A_\mu = \partial_\mu \lambda \), \( \delta \sigma = -m \lambda \). With the gauge fixing term \( \mathcal{L}_{gf} = -\left( \partial_\mu A^\mu + \xi m \sigma \right)^2 / 2 \xi \), the total Lagrangian reads

\[
\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{m^2}{2} A_\mu A^\mu - \frac{1}{2 \xi} (\partial_\mu A^\mu)^2 - \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma - \xi \frac{m^2}{2} \sigma^2 + g A_\mu J^\mu,
\]

where we have added also an interaction term which contains the coupling of \( A_\mu \) with fermions via a conserved current with \( \partial_\mu J^\mu = 0 \). Here the fields \( \sigma \) and \( A_\mu \) are decoupled and renormalizability and unitarity are manifest. Mass growth by the Stueckelberg mechanism occur quite naturally D brane constructions where one encounters the group \( U(N) \) for a stack of \( N \) D branes which is then broken to its subgroup \( SU(N) \) via
Stueckelberg couplings. Thus, for example, one has

\[
\mathcal{L}_{\text{St}} = -\frac{1}{4} C_{\mu} C_{\nu} + g_{X} C_{\mu} J_{X}^\mu - \frac{1}{2}(\partial_{\mu}\sigma + M_{1} C_{\mu} + M_{2} B_{\mu})^2.
\]  

(4)

It is easily checked that the above Lagrangian is invariant under the following transformations: \( \delta_{Y}(C_{\mu}, B_{\mu}, \sigma) = (0, \partial_{\mu} \lambda_{Y}, -M_{2} \lambda_{Y}) \) and \( \delta_{X}(C_{\mu}, B_{\mu}, \sigma) = (\partial_{\mu} \lambda_{X}, 0, -M_{1} \lambda_{X}) \).

The two Abelian gauge bosons can be decoupled from \( \sigma \) by the addition of gauge fixing terms as before. Additionally, of course, one has to add the standard gauge fixing terms for the SM gauge bosons to decouple from the Higgs.

We look now at the physical content of the theory. In the vector boson sector in the basis \( V_{\mu}^{T} = (C_{\mu}, B_{\mu}, A_{\mu}^{3}) \), the mass matrix for the vector bosons takes the form

\[
M_{V}^2 = \begin{bmatrix}
M_{1}^2 & M_{1} M_{2} & 0 \\
M_{1} M_{2} & M_{2}^2 + \frac{1}{4} v^2 g_{Y}^2 & -\frac{1}{2} v^2 g_{2} g_{Y} \\
0 & -\frac{1}{4} v^2 g_{2} g_{Y} & \frac{1}{4} v^2 g_{2}^2
\end{bmatrix},
\]  

(5)

where \( g_{2} \) and \( g_{Y} \) are the \( SU(2)_{L} \) and \( U(1)_{Y} \) gauge coupling constants, and are normalized so that \( M_{W}^2 = g_{2}^2 v^2 / 4 \). It is easily checked that \( \det(M_{V}^2) = 0 \) which implies that one of the eigenvalues is zero, whose eigenvector we identify with the photon. The remaining two eigenvalues are non-vanishing and correspond to the \( Z \) and \( Z' \) bosons. The symmetric matrix \( M_{V}^2 \) can be diagonalized by an orthogonal transformation, \( V = \Theta E \), with \( E_{\mu}^{T} = (Z_{\mu}, Z_{\mu}, A_{\mu}^{3}) \) so that the eigenvalues are given by the set: \( \text{diag}(M_{V}^2) = \{ M_{Z}^2, M_{Z'}^2, 0 \} \).

One can solve for \( \Theta \) explicitly and we use the parametrization

\[
\Theta = \begin{bmatrix}
\cos \psi \cos \phi - \sin \theta \sin \phi \sin \psi & -\sin \psi \cos \phi - \sin \theta \sin \phi \cos \psi & -\cos \theta \sin \phi \\
\sin \psi \sin \phi + \sin \theta \cos \phi \sin \psi & -\sin \psi \sin \phi + \sin \theta \cos \phi \cos \psi & \cos \theta \cos \phi \\
-\cos \theta \sin \psi & -\cos \theta \cos \psi & \sin \theta
\end{bmatrix},
\]

where \( \tan(\phi) = M_{Z} / M_{1} \equiv \epsilon \), \( \tan(\theta) = g_{2} / g_{Y} \cos(\phi) = \tan(\theta_{W}) \cos(\phi) \). The third angle is given by \( \tan(\psi) = \tan(\theta) \tan(\phi) M_{W}^2 / (\cos(\theta) (M_{Z}^2 - M_{W}^2(1 + \tan^2(\theta)))) \). This allows one to choose \( \epsilon \) and \( M_{1} \) as two independent parameters to characterize physics beyond SM. There is also a modification of the expression of the electric charge in terms of SM parameters. Thus if we write the EM interaction in the form \( e A_{\mu}^{0} J_{\text{em}}^{\mu} \) the expression for \( e \) is given by

\[
e = g_{2} g_{Y} \cos(\phi) / \sqrt{g_{2}^2 + g_{Y}^2 \cos^2(\phi)}.
\]  

(6)
The LEP and Tevatron data puts stringent bounds on \( \epsilon \). One finds [9, 10] that it is constrained by \( \epsilon \lesssim 0.06 \) in most of the parameter space. In the absence of a hidden sector, i.e., the matter sector that couples only to \( C_\mu \), the \( Z' \) can decay only into visible sector quarks and leptons, and its decay width is governed by \( \epsilon \) and hence the \( Z' \) is very sharp, with a width that lies in the range of of maximally several hundred MeV compared to several GeV that one expects for a \( Z' \) arising from a GUT group (a narrow \( Z' \) can also arise in other models, see e.g., [11, 12, 13, 14, 15]). However, even a very sharp \( Z' \) is discernible at the Tevatron and at the LHC using the dilepton signal. On the other hand, in the absence of a hidden sector, \( \hat{\epsilon} \) can decay only into visible sector particles and will have a width in the several GeV range. In this case the branching ratio of \( Z' \) to \( l^+l^- \) will be very small [16, 17] and the dilepton signal will not be detectable. We will return to this issue in the context of milli charged dark matter.

**STUECKELBERG EXTENSION OF THE MINIMAL SUPERSYMMETRIC STANDARD MODEL**

To obtain the supersymmetric Stueckelberg extension [6, 8, 7] we consider the Stueckelberg chiral multiplet \( S = (\rho + i\sigma, \chi, F_S) \) along with the vector superfield multiplets for the \( U(1)_Y \) denoted by \( B = (B_\mu, \lambda_B, D_B) \) and for the \( U(1)_X \) denoted by \( C = (C_\mu, \lambda_C, D_C) \). The Stueckelberg addition to the SM Lagrangian is then given by

\[
\mathcal{L}_{St} = \frac{1}{2} d^2 \theta d^2 \bar{\theta} \ (M_1 C + M_2 B + S + \bar{S})^2. \tag{7}
\]

Under \( U(1)_Y \) and \( U(1)_X \) the supersymmetrized gauge transformations are then given by:

\[
\delta_Y (C, B, S) = \{0, \Lambda_Y + \tilde{\Lambda}_Y, -M_2 \Lambda_Y \} \quad \text{and} \quad \delta_X (C, B, S) = \{\Lambda_X + \tilde{\Lambda}_X, 0, -M_1 \Lambda_X \}.
\]

Expanding the fields in the component form, in the Wess-Zumino gauge, we have for a vector superfield, denoted here by \( V = (C, B) \),

\[
V = -\theta \sigma^\mu \bar{\theta} V_\mu + i\theta \bar{\theta} \bar{\theta} \lambda^\mu - i\bar{\theta} \bar{\theta} \theta \lambda^\mu + \frac{1}{2} \theta \theta \bar{\theta} \bar{\theta} D_V. \tag{8}
\]

The superfield \( S \) in component notation is given by

\[
S = \frac{1}{2} (\rho + i\sigma) + \theta \chi + i\theta \sigma^\mu \bar{\theta} \frac{1}{2} (\partial_\mu \rho + i\partial_\mu \sigma) + \theta \theta F_S + \frac{i}{2} \theta \theta \bar{\theta} \sigma^\mu \partial_\mu \chi + \frac{1}{8} \theta \theta \bar{\theta} \bar{\theta} (\Box \rho + i\Box \sigma). \tag{9}
\]

We note that the superfield \( S \) contains the scalar \( \rho \) and the axionic pseudo-scalar \( \sigma \). In component form \( \mathcal{L}_{St} \) then has the form

\[
\mathcal{L}_{St} = \frac{1}{2} (M_1 C_\mu + M_2 B_\mu + \partial_\mu \sigma)^2 - \frac{1}{2} (\partial_\mu \rho)^2 - i\chi \sigma^\mu \partial_\mu \bar{\chi} + 2 |F_S|^2 \tag{10}
\]

\[+\rho (M_1 \lambda_C + M_2 \lambda_B) + [\chi (M_1 \lambda_C + M_2 \lambda_B) + \text{h.c.}].\]
To the above we can add the gauge fields of the Standard Model which give

\[ \mathcal{L}_{\text{gauge}} = -\frac{1}{4} C_{\mu \nu} C^{\mu \nu} - \frac{1}{4} B_{\mu \nu} B^{\mu \nu} - i \lambda_B \sigma^\mu \partial_\mu \tilde{\lambda}_B - i \lambda_C \sigma^\mu \partial_\mu \tilde{\lambda}_C + \frac{1}{2} D^2_C + \frac{1}{2} D^2_B. \]

The gauge fields can be coupled to the chiral superfields \( \Phi_i \) of matter in the usual way

\[ \mathcal{L}_{\text{matter}} = \int d^2 \theta d^2 \bar{\theta} \left[ \sum_i \Phi_i e^{2gY} Q_Y B + 2gX Q_X C \Phi_i + \sum_i \bar{\Phi}_{\text{hid},i} e^{2gY} Q_Y B + 2gX Q_X C \Phi_{\text{hid},i} \right]. \]

Here \( Q_Y = Y/2 \), and where \( Y \) is the hypercharge so that \( Q = T_3 + Y/2 \). We assume that the SM matter fields do not carry any charge under the hidden gauge group, i.e. \( Q_X \Phi_i = 0 \). The Stueckelberg extensions of the type we have discussed could have origin in Type IIB orientifold models [18, 19, 20, 21, 22] and several recent works appear to recover in its low energy limit the type of models discussed here [23, 24, 25, 26, 27, 28, 29].

**Milli weak dark matter in \( U(1)_X \) extension**

We note that the Stueckelberg extension brings in two more Majorana spinors which we can construct out of the Weyl spinors as follows \( \psi_0^T = (\chi^\alpha, \tilde{\chi}^\dot{\alpha}) \), \( \lambda^T = (\lambda_\alpha, \tilde{\lambda}^\dot{\alpha}) \). This enlarges the neutralino mass matrix from being 4 × 4 as is the case in MSSM to a 6 × 6 mass matrix in the Stueckelberg extension. The enlarged neutralino mass matrix reads

\[
M_{1/2} = \begin{bmatrix}
0 & M_1 & M_2 & 0 & 0 & 0 \\
M_1 & \tilde{m}_X & 0 & 0 & 0 & 0 \\
M_2 & 0 & \tilde{m}_1 & 0 & -c_{\beta} s_{SW} M_0 & s_{\beta} s_{SW} M_0 \\
0 & 0 & 0 & \tilde{m}_2 & c_{\beta} c_{CW} M_0 & -s_{\beta} c_{CW} M_0 \\
0 & 0 & -c_{\beta} s_{SW} M_0 & c_{\beta} c_{CW} M_0 & 0 & -\mu \\
0 & 0 & s_{\beta} s_{SW} M_0 & -s_{\beta} c_{CW} M_0 & -\mu & 0
\end{bmatrix}. \tag{11}
\]

Here the 4 × 4 matrix on the lower right hand corner is the usual neutralino mass matrix of MSSM, while the 2 × 2 matrix in the top left hand corner is due the Stueckelberg extension. The term \( \tilde{m}_X \) is the soft breaking term which is added by hand. The zero entry in the upper left hand corner arises due to the Weyl fermions not acquiring soft masses. The 6 × 6 matrix gives rise to six Majorana mass eigenstates which may be labeled as follows \( E_{[1/2]} = (\chi^0_1, \chi^0_2, \chi^0_3, \chi^0_4, \chi^0_5, \chi^0_6)^T \), where the two additional Majorana eigenstates \( (\chi^0_5, \chi^0_6) \) are due to the Stueckelberg extension. We label these two \( \xi^0_1, \xi^0_2 \) and to leading order in \( \epsilon \) their masses are given by

\[
m_{\xi^0_1} \simeq \sqrt{M^2 + \frac{1}{4} \tilde{m}_X^2 - \frac{1}{2} \tilde{m}_X}, \quad m_{\xi^0_2} \simeq \sqrt{M^2 + \frac{1}{4} \tilde{m}_X^2 + \frac{1}{2} \tilde{m}_X}. \tag{12}
\]

where \( M^2 = M^2_1 + M^2_2 \). If the mass of \( \xi^0_1 \) is less than the mass of other sparticles, then \( \xi^0_1 \) will be a candidate for dark matter with R parity conservation. These are what one may
call XWIMPS (mWIMPS) for extra (milli) weakly interacting massive particles. Here the satisfaction of relic density requires coannihilation and one has to consider processes of the type $\xi^0 + \xi^0 \rightarrow X$, $\xi^0 + \chi^0 \rightarrow X'$, $\chi^0 + \chi^0 \rightarrow X''$, where \{X\} etc denote the Standard Model final states. In this case we can write the effective cross section as follows \cite{30}

$$\sigma_{\text{eff}} = \frac{1}{\sigma_{\xi^0 \chi^0}} \frac{1}{\sigma_{\chi^0 \chi^0}} \left( Q + \frac{\sigma_{\xi^0 \chi^0}}{\sigma_{\chi^0 \chi^0}} \right)^2, \quad Q = \frac{g_{\chi^0}}{g_{\xi^0}} (1 + \Delta) e^{-x_f \Delta}. \quad (13)$$

Here $g$ is the degeneracy for the corresponding particle, $x_f = m_{\xi^0}/T_f$ where $T_f$ is the freeze-out temperature, and $\Delta = (m_{\chi^0} - m_{\xi^0})/m_{\xi^0}$ is the mass gap. For the case of XWIMPS one has $\sigma_{\xi^0 \chi^0}/\sigma_{\chi^0 \chi^0} \sim O(\epsilon^2) \ll 1$. Now it is easily seen that when the mass gap between $\xi^0$ and $\chi^0$ is large and $x_f \Delta \gg 1$, then $\sigma_{\text{eff}}$ is much smaller than the typical WIMP cross-section and in this case one does not have an efficient annihilation of the XWIMPS. On the other hand if the mass gap between the XWIMP and WIMP is small then coannihilation of XWIMPs is efficient. In this case $Q \sim 1$ and one has $\sigma_{\text{eff}} \simeq \sigma_{\xi^0 \chi^0} \left( \frac{Q}{1+Q} \right)^2$. The above result is valid more generally with many channels participating in the coannihilations, as can be seen by defining an effective $Q$ given by $Q = \sum_{i=2}^N Q_i$ where $Q_i = (g_i/g_1)(1 + \Delta_i)^{3/2} e^{-x_f \Delta_i}$. Thus, satisfaction of the relic density constraints arise quite easily for the XWIMPS. A detailed analysis of the relic density of XWIMPS was carried out in \cite{30} and it was found that the WMAP-3 constraint $\Omega_{CDM} h^2 = 0.1045_{+0.0072}^{-0.0095}$ can be satisfied by XWIMPS.

**STUECKELBERG MECHANISM WITH KINETIC MIXING**

We discuss now the Stueckelberg extension with kinetic mixing \cite{17} for which we take the Lagrangian to be of the form $\mathcal{L}_{\text{StkSM}} = \mathcal{L}_{\text{SM}} + \Delta \mathcal{L}$ where

$$\Delta \mathcal{L} \supset \frac{1}{4} C_{\mu \nu} C^{\mu \nu} - \frac{\delta}{2} C_{\mu \nu} B^{\mu \nu} - \frac{1}{2} (\partial_{\mu} \sigma + M_1 C_{\mu} + M_2 B_{\mu})^2 + g_X J^\mu_X \sigma \mu. \quad (14)$$

In this case the kinetic mixing matrix, in the basis $V^T = (C, B, A^3)$ is,

$$\mathcal{K} = \begin{bmatrix} 1 & \delta & 0 \\ \delta & 1 & 0 \\ 0 & 0 & 1 \end{bmatrix}. \quad (15)$$

A simultaneous diagonalization of the kinetic energy and of the mass matrix can be obtained by a transformation $T = KR$, which is a combination of a $GL(3)$ transformation ($K$) and an orthogonal transformation ($R$). This allows one to work in the diagonal basis, denoted by $E^T = (Z', Z, A')$, through the transformation $V = (KR) E$, where the matrix $K$ which diagonalizes the kinetic terms has the form

$$K = \begin{bmatrix} C_\delta & 0 & 0 \\ -S_\delta & 1 & 0 \\ 0 & 0 & 1 \end{bmatrix}, \quad C_\delta = \frac{1}{\sqrt{1 - \delta^2}}, \quad S_\delta = \delta C_\delta. \quad (16)$$
The colored regions indicate the satisfaction of the relic density constraints consistent with the WMAP-3 constraints and the size of the dilepton signal \( \sigma \cdot Br(Z' \rightarrow l^+ l^-) \) at the Tevatron as a function of \( M_{Z'} \) when \( 2M_{Z'} = 300 \text{ GeV} \). The curves in ascending order are for values of \( \bar{\epsilon} \) in the range \((0.01 - 0.06)\) in steps of 0.01. The dilepton signal has a dramatic fall as \( M_{Z'} \) crosses the point \( 2M_{Z'} = 300 \text{ GeV} \) where the \( Z' \) decay into the hidden sector fermions is kinematically allowed, widening enormously the \( Z' \) decay width. The green shaded regions are where the WMAP-3 relic density constraints are satisfied for the case when there is no kinetic mixing. Red and blue regions are for the case when kinetic mixing is included. The current constraints on the dilepton and signal from CDF\[32\] and the DØ search for narrow resonances \[33\] are also exhibited. From \[17\].

The diagonalization also leads to the following relation for the electronic charge

\[
\frac{1}{e^2} = \frac{1}{g_2^2} + \frac{1-2\bar{\epsilon}\delta + \bar{\epsilon}^2}{g_Y^2}.
\]

(17)

Thus \( g_Y \) is related to \( g_Y^{SM} \) by

\[
g_Y = \gamma \sqrt{1 + \bar{\epsilon}^2 - 2\bar{\epsilon}\delta}, \quad \gamma \equiv g_Y^{SM}.
\]

In the absence of a hidden sector, there is only one parameter that enters in the analysis of electroweak fits. This effective parameter is given by \( \bar{\epsilon} = (\epsilon - \delta) / \sqrt{1 - \delta^2} \). Thus one can satisfy the LEP and the Tevatron electro-weak data with \( \bar{\epsilon} \lesssim 0.06 \) but \( \epsilon \) and \( \delta \) could be individually larger.

**How milli charge is generated in Stueckelberg extension**

To exhibit the phenomenon of generation of milli-charge in the Stueckelberg model we consider two gauge fields \( A_{1\mu}, A_{2\mu} \) corresponding to the gauge groups \( U(1) \) and \( U(1)' \). We choose the following Lagrangian \( \mathcal{L} = \mathcal{L}_0 + \mathcal{L}_1 + \mathcal{L}_2 \) where

\[
\mathcal{L}_0 = -\frac{1}{4} F_{1\mu\nu} F_1^{\mu\nu} - \frac{1}{4} F_{2\mu\nu} F_2^{\mu\nu} - \frac{\delta}{2} F_{1\mu\nu} F_2^{\mu\nu}, \quad \mathcal{L}_1 = J_1^\mu A_1^\mu + J_\mu A_2^\mu,
\]

\[
\mathcal{L}_2 = -\frac{1}{2} M_1^2 A_1^\mu A_1^\mu - \frac{1}{2} M_2^2 A_2\mu A_2^\mu - M_1 M_2 A_{1\mu} A_{2}^\mu.
\]

(18)
Here $J_{\mu}$ is the current arising from the physical sector including quarks, leptons, and the Higgs fields and $J'_{\mu}$ is the current arising from the hidden sector. As indicated in the discussion preceding Eq. (16), the mass matrix can be diagonalized by the $R$ transformation which for this $2 \times 2$ example is parameterized as follows

$$
R = \begin{bmatrix}
\cos \theta & -\sin \theta \\
\sin \theta & \cos \theta
\end{bmatrix},
$$

where $\theta$ is determined by the diagonalization constraint so that

$$
\theta = \arctan \left( \frac{\epsilon \sqrt{1 - \delta^2}}{1 - \delta \epsilon} \right). \tag{20}
$$

The diagonalization yields one massless mode $A^\mu_\gamma$ and one massive mode $A^\mu_M$. In this case the interaction Lagrangian in the diagonal basis assumes the form

$$
\mathcal{L}_1 = \frac{1}{\sqrt{1 - 2 \delta \epsilon + \epsilon^2}} \left( \frac{\epsilon - \delta}{\sqrt{1 - \delta^2}} J_{\mu} + \frac{1 - \delta \epsilon}{\sqrt{1 - \delta^2}} J'_{\mu} \right) A^\mu_M + \frac{1}{\sqrt{1 - 2 \delta \epsilon + \epsilon^2}} \left( J_{\mu} - \epsilon J'_{\mu} \right) A^\mu_\gamma. \tag{21}
$$

The interesting phenomenon to note here is that the photon field $A^\mu_\gamma$ couples with the hidden sector current $J'_{\mu}$ only due to mass mixing, i.e., only due to $\epsilon$. Thus the origin of milli charge is due to the Stueckelberg mass mixing both in the presence or absence of kinetic mixing. This phenomenon persists when one considers $G_{SM} \times U(1)_X$ where the $SU(2)_L \times U(1)_Y$ gauge group is broken by the conventional Higgs mechanism and in addition one has the Stueckelberg mechanism generating a mass mixing between the $U(1)_Y$ and $U(1)_X$. The above phenomenon is to be contrasted with the kinetic mixing model [34] where one has two massless modes (the photon and the paraphoton) and the photon can couple with the hidden sector because of kinetic mixing generating milli charge couplings. [An analysis with kinetic mixing and mass mixings of a different type than discussed here is also considered in [35]].

**Milli charge dark matter**

The hidden sector particles are typically natural candidates for dark matter. The main issue concerns their ability to annihilate in sufficient amounts to satisfy the current relic density constraints. Now the milli charged particles could decay in sufficient amounts by decaying via the $Z'$ to the Standard Model particles if their masses are $< M_{Z'}/2$. An explicit analysis of this possibility is carried out in [16] where a pair of Dirac fermions were put in the hidden sector which couple with strength $g_2$ with the Stueckelberg field $C_{\mu}$. In this case it was shown that the relic density constraints consistent with the WMAP-3 data can be satisfied. Further, with inclusion of proper thermal averaging of the quantity $\langle \sigma v \rangle$ over the resonant $Z'$ [using techniques discussed in [36, 37, 38, 39, 40]] which enters in
FIGURE 2. An analysis of the relic density of milli-charged particles for the case when kinetic mixing is included in the Stueckelberg $Z'$ model. The analysis is done for $M_Z = 150$ GeV, $\bar{\epsilon} = 0.04$, and $\delta = (0.05, 0.075, 0.10, 0.15, 0.20, 0.25)$, where the values are in descending order for $M_{Z'} > 300$ GeV. The red and black bands are the WMAP-3 constraints where the black band also produces an observable dilepton signal. The analysis shows that for $\bar{\epsilon}$ fixed, increasing $\delta$ increases the parameter space where the WMAP-3 relic density constraint is satisfied, while allowing for a detectable $Z$ prime signal as shown in Fig. (1). From [17].

the relic density analysis, one finds that the WMAP-3 relic density constraints can also be satisfied over a broad range when the masses of the milli charged hidden sector particles lie above $M_{Z'}/2$, with and without kinetic mixing[17]. This phenomenon comes about because of the thermal averaging effect. On the branch where the milli charged particles have masses lying above $M_{Z'}/2$ the relic density constraints can be satisfied and still produce a dilepton signal which may be observable at the LHC. [17]. Satisfaction of the relic density constraints consistent with WMAP-3 and illustration of the strong dilepton signal are seen in Figs. (12)[taken from [17]]. The experimental constraints on milli charged particles have been discussed in a number of papers in the literature mostly in the context of kinetic mixing models, [41, 42, 43, 44, 45, 46, 47, 48, 49, 50, 51, 52, 53], but without mass generation via the Stueckelberg mechanism.

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