Atomki Anomaly in Family-Dependent $U(1)'$ Extension of the Standard Model

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In the context of a gauge invariant, nonanomalous, and family-dependent (nonuniversal) $U(1)'$ extension of the Standard Model, wherein a new high-scale mechanism generates masses and couplings for the first two fermion generations and the standard Higgs mechanism does so for the third one, we find solutions to the anomaly observed by the Atomki Collaboration in the decay of excited states of beryllium, in the form of a very light $Z'$ state, stemming from the $U(1)'$ symmetry breaking, with significant axial couplings so as to evade a variety of low-scale experimental constraints.

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I. INTRODUCTION

The Atomki Collaboration [1] has recently detected hints of a new light bosonic state, with mass $\approx 17$ MeV, from the measurement of the angle between $e^+e^-$ pairs and their invariant mass produced by the 18.15 MeV nuclear transition in the excited state $^8\text{Be}^*$ [2] (see also Refs. [3–6]).1 There have been several studies [7–19] trying to explain the nature of this new state which mostly focus on a vector boson solution. In this work, we further consider this possible scenario in the context of a rather minimal model: specifically, by extending the Standard Model (SM) with a single family-dependent (nonuniversal) $U(1)'$ group.

As the model contains two Abelian groups, $U(1)_L \times U(1)'$, there will be a mixing between the hypercharge gauge boson $B_\mu$ of the SM and the new $U(1)'$ gauge boson $B'_\mu$. Therefore, the kinetic Lagrangian is given by

$$\mathcal{L}_{\text{kin}} = -\frac{1}{4} \tilde{F}_{\mu\nu} \tilde{F}^{\mu\nu} - \frac{1}{4} \tilde{F}'_{\mu\nu} \tilde{F}'^{\mu\nu} - \frac{\kappa}{2} \tilde{F}_\mu \tilde{F}'^\mu, \quad (1)$$

where $\kappa$ parametrizes the level of mixing between the two fields. One may diagonalize the kinetic Lagrangian by a rotation and rescaling of these fields, which leaves the covariant derivative as

$$\mathcal{D}_\mu = \partial_\mu + \cdots + ig_1 Y B_\mu + i(\tilde{g} Y + g' z) B'_\mu, \quad (2)$$

where $Y$ and $g_1$ are the hypercharge and its gauge coupling, $z$ and $g'$ are the $U(1)'$ charge and its gauge coupling, and $\tilde{g}$ is the mixed gauge coupling between the groups. We break the $U(1)'$ with a new SM-singlet scalar, $\chi$, with a charge $z_\chi$ under the new gauge group, with a vacuum expectation value (VEV) $\langle \chi \rangle = v'$ inducing a mass term $m_{\chi} = g' z_\chi v'$. It is interesting to note that, for $g' \sim \mathcal{O}(10^{-4} - 10^{-5})$, as required by several experimental constraints, $m_{\chi}$ can be of order $\mathcal{O}(10)$ MeV if $v'$, the scale of $U(1)'$ symmetry breaking, is of order $\mathcal{O}(100–1000)$ GeV.

This massive vector boson interacts with the SM fermions through the gauge current

$$J_{\mu}^f = \sum_{f} \bar{\psi}_f J_{\mu}^{f}(C_{f,L} P_L + C_{f,R} P_R) \psi_f, \quad (3)$$

with left- ($L$) and right- ($R$) handed coefficients [19]

$$C_{f,L} = -g_Z s_f' (T_f' - s_W Q_f') + (\tilde{g} Y + g' z_{f,L}) c_f',$$

$$C_{f,R} = g_Z s_w^2 Q_f + (\tilde{g} Y_{f,R} + g' z_{f,R}) c_f', \quad (4)$$

where we have defined $g_Z = \sqrt{g_1^2 + g'^2}$ [the electroweak (EW) coupling], $s_w \equiv \sin(\theta_W)$, $c_w \equiv \cos(\theta_W)$, $s' \equiv \sin(\theta')$, and $c' \equiv \cos(\theta')$, with $\theta'$ being the angle parametrizing the

1In fact, the 17.64 MeV transition also eventually appeared to present a similar anomaly, albeit less significant, with a boson mass broadly compatible with the previous one; however, it should be mentioned that this was never documented in a published paper, only in proceeding contributions, so we do not consider it here.

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aforementioned gauge mixing, and we have also introduced $T^3_f$ and $Q_f$, the weak isospin and electric charge of the fermion $f$, respectively. Finally, $Y_{f,L/R}$ and $z_{f,L/R}$ represent the hypercharge and $U(1)'$ quantum numbers of the $L$- or $R$-handed fermion. By diagonalizing the mass matrix of neutral gauge bosons, one finds this mixing angle, $\theta'$, effectively between the SM $Z$ and the new $Z'$ [associated with $U(1)'$], as [20]

$$\tan 2\theta' = \frac{2g_H g_Z}{g_H^2 + 4m_W^2/v^2 - g_Z^2},$$  
(5)

where $g_H = \tilde{g} + 2g'z_H$.

We now define the usual vector ($V$) and axial ($A$) coefficients in the limit of small gauge coupling and mixing, $g', \tilde{g} \ll 1$:

$$C_{f,V} = \frac{C_{f,R} + C_{f,L}}{2} \approx g'z_H(T^3_f - 2z_2^2 Q_f) + z_{f,V},$$  
(6)

$$C_{f,A} = \frac{C_{f,R} - C_{f,L}}{2} \approx g'[-z_H T^3_f + z_{f,A}],$$  
(7)

where we use the convention $Y_f = Q_f - T^3_f$ and define the $V$ and $A$ quantum numbers under the $U(1)'$ group, $z_{f,V/A} = 1/2(z_{f,R} \pm z_{f,L})$.

The Yukawa sector of the SM for quarks and leptons takes the form

$$-\mathcal{L}_{\text{Yuk}} = Y_u \bar{Q} H u_R + Y_d \bar{Q} H d_R + Y_e \bar{L} H e_R.$$  
(8)

Because of gauge invariance, this imposes a condition on the combination of charges of the fields under the $U(1)'$ group:

$$-z_u - z_H + z_d = -z_Q + z_H + z_e = 0.$$  
(9)

Inserting these relations into Eq. (7), one finds no $A$ couplings to the $Z'$ for quarks and leptons; i.e., $C_{(q,f)^+}, A \approx 0$ at leading order in the gauge coupling $g'$.

It is difficult to construct a model (with minimal extra particle content) with only V interactions of fermions to the $Z'$, as opposed to $A$, because relatively larger couplings$^2$ are required to achieve a successfully high rate for the transition $^8\text{Be} \rightarrow ^7\text{BeZ}'$, possibly explaining the Atomki anomaly. This is because the contributions of $A$ couplings in the transition are proportional to $k/M_{Z'}^2 \ll 1$ (where $k$ is the small momentum of the $Z'$), whereas the $V$ component has a momentum proportionality of $k^3/M_{Z'}^3$, as explained in Ref. [8].

In the (purely) $V$ case, the larger values of $(g, \tilde{g})$ conflict with the nonobservation of deviations from the SM by neutrino scattering off electrons (see below—in fact, we detail these experimental requirements on our particular model construction later on). One possibility, explored in Ref. [19], is to employ a two-Higgs doublet model (2HDM), which successfully augments the Yukawa sector such that this condition of gauge invariance is modified by the second Higgs doublet and eventually allows for nonsuppressed $A$ couplings. This ensures that the Atomki anomaly can be explained with smaller $g', \tilde{g}$ gauge couplings, which thus alleviate the present experimental constraints.

In this work, we proceed in a different direction. Namely, to allow for $A$ couplings, we consider the possibility of having a family-dependent (nonuniversal) $U(1)'$. In this case, the Yukawa interaction terms, in Eq. (8), are modified as follows:

$$-\mathcal{L}_{\text{Yuk}} = \Gamma_u \frac{z_{Q_i}}{M_{\mu^{Q_i}}} \bar{Q}_L iH u_{R,j} + \Gamma_d \frac{z_{Q_i}}{M_{\mu^{Q_i}}} \bar{Q}_L iH d_{R,j} + \Gamma_e \frac{z_{e_i}}{M_{\mu^{e_i}}} \bar{L}_i H e_{R,j} + \text{H.c.},$$  
(10)

where the dimension of the nonrenormalizable scale $M$ is specified by the $U(1)'$ charges of the involved fields. This procedure can be used to generate fermion masses at tree level or at higher orders [21].$^3$ Therefore, here we assume $U(1)'$ charges such that the third fermion family Yukawa structure is SM-like, due to more natural $O(1)$ couplings, while the masses of the first two quark and lepton families can be obtained through some higher-order corrections. In fact, various models attempt to explain the smallness of the first two quark and lepton families by a radiative mass generation mechanism, as in Ref. [22], or by horizontal symmetries [21]. Explicitly, we require that the condition in Eq. (9) only hold for the third generation. In short, we choose to impose that the first two generations be flavor universal, but not the third. $z_{i} = z_{i_L}$ for $i = \{Q, u_R, d_R, L, e_R\}$.

In addition to the aforementioned conditions of gauge invariance of the third-generation Yukawa couplings and flavor universality in the first two generations, we now discuss some additional constraints on our charge assignment. Despite working with a low-scale, phenomenological approach, we choose to adhere to the chiral anomaly cancellation conditions satisfied by the current fermionic content of the SM in addition to $R$-handed neutrinos. The six anomaly conditions are summarized as

$$\sum_i^3 (2z_{Q_i} - z_{u_i} - z_{d_i}) = 0,$$  
(11)

$^2$Though still in the regime $(g, \tilde{g}) \ll 1$.

$^3$It may be interesting to investigate whether the same $U(1)'$ symmetry that explains the Atomki anomaly could act as a flavor symmetry and arrange for the observed fermion mass hierarchy and mixing. However, this is beyond the scope of this paper.
We thus impose that there must be no couplings at all for neutrinos to the TEXONO experiment [7,24], and by the electron-neutrino scattering by the TEXONO experiment [23], and by the electron-neutrino scattering by the TEXONO experiment [7,24–26]. We thus impose that there must be no couplings at all for neutrinos to the $Z'$: $C_{v,A} = C_{v,L} = 0$. One finds then the additional requirement that

$$Z_{d_i} = Z_{L_3} = -Z_H.$$  \[ (17) \]

As stated before, we also require $A$ couplings for the first two generations of quarks to successfully reproduce the Atomki anomaly:

$$-Z_{Q_{1,2}} - Z_H + Z_{d_{1,2}} \neq 0,$$  \[ (18) \]

$$-Z_{Q_{1,2}} + Z_H + Z_{d_{1,2}} \neq 0.$$  \[ (19) \]

However, $A$ couplings in the charged lepton sector have stringent constraints from atomic parity violation in cesium (Cs) [27]. These can be extracted from the measurement of the effective weak charge $\Delta Q_W$ of the Cs atom:

$$\Delta Q_W = -\frac{\sqrt{2}}{G_F} C_{e,A}(C_{u,V}(2Z + N) + C_{d,V}(Z + 2N)) \times \left(\frac{0.8}{(17 \text{ MeV})^2}\right) \lesssim 0.71.$$  \[ (20) \]

As the vector couplings of the $Z'$ to up and down quarks are, in general, nonzero, we thus also require that there be no $A$ interaction with the electrons:

$$C_{e,A} = 0.$$  \[ (21) \]

This will also help to alleviate bounds from $(g - 2)_e$ which are discussed later. For the same reason, we also set the muon $A$ coupling to zero, to avoid increasing the discrepancy between the experimental measurement and the SM prediction of the $(g - 2)_\mu$ (discussed further in the paper),

$$C_{\mu,A} = 0.$$  \[ (22) \]

With these final constraints, we find that our initial 16 free charges (three generations of $\{Z_{Q_i}, Z_{u_i}, Z_{d_i}, Z_{L_i}, Z_{e_i}\}$ and $Z_H$) may be expressed as a function of one single parameter. Adjusting this parameter is equivalent to a rescaling of the coupling, so our charge assignment with these constraints is fixed (see Table I), and we normalize it with $Z_H = 1$.

### II. CONSTRAINTS

We now discuss the Atomki anomaly requirements and the experimental constraints quantitatively.

The Atomki Collaboration [2] has published that the best fit for the mass of the (would-be) $Z'$ should be $M_{Z'} = 16.7 \pm 0.35(\text{stat}) \pm 0.5(\text{sys}) \text{MeV}$, corresponding to a ratio of branching ratios (BRs),

$$\text{BR}(8^7\text{Be}^* \to Z' + 8^7\text{Be}) \times \text{BR}(Z' \to e^+ e^-) = 5.8 \times 10^{-6},$$  \[ (23) \]

with a statistical significance of $\sim 6\sigma$ [2]. However, the Atomki Collaboration has since then pursued further masses and BRs, as mentioned in Ref. [7], as a private

\[ TABLE I. Charge assignment of the SM particles under the family-dependent (nonuniversal) $U(1)'$. This numerical charge assignment satisfies the discussed anomaly cancellation conditions, and it enforces a gauge invariant Yukawa sector of the third generation and family universality in the first two fermion generations, as well as no coupling of the $Z'$ to the all-neutrino generations. \]

| $SU(3)$ | $SU(2)$ | $U(1)_Y$ | $U(1)'$ |
|--------|--------|---------|---------|
| $Q_1$  | 3      | 2       | 1/6     |
| $Q_2$  | 3      | 2       | 1/6     |
| $Q_3$  | 3      | 2       | 1/6     |
| $u_R$  | 3      | 1       | 2/3     |
| $d_R$  | 3      | 1       | 2/3     |
| $e_L$  | 3      | 1       | 2/3     |
| $e_R$  | 3      | 1       | 2/3     |
| $L_1$  | 1      | 2       | -1/2    |
| $L_2$  | 1      | 2       | -1/2    |
| $L_3$  | 1      | 2       | -1/2    |
| $H$    | 1      | 2       | 1/2     |

\[ PHYS. REV. D 99, 055022 (2019) 055022-3 \]
Before evaluating the region of the parameter space explaining the anomalous $^8\text{Be}$ transition, though, we ought to discuss in more detail the various experimental constraints which affect such a low-mass and weakly coupled $Z'$. First, we have not seen such a $Z'$ in electron beam dump experiments (e.g., SLAC E141) \[31,32\]. Therefore, the $Z'$ has not been produced herein, hence

$$C_{e,V}^2 + C_{e,A}^2 < 10^{-17},$$

or else the $Z'$ has been caught in the dump, hence

$$\frac{C_{e,V}^2 + C_{e,A}^2}{\text{BR}(Z' \to e^+e^-)} \gtrsim 3.7 \times 10^{-9}. \quad (33)$$

As the former is not compatible with the Atomki observation, we will consider the latter condition. We have also not seen the $Z'$ in the NA64 beam dump experiment \[33\], which places the (stronger than E141) bound,

$$\frac{C_{e,V}^2 + C_{e,A}^2}{\text{BR}(Z' \to e^+e^-)} \gtrsim 1.6 \times 10^{-8}. \quad (34)$$

We have not seen a $Z'$ in parity-violating Moller scattering (e.g., SLAC E158) \[34\]. Therefore, the following constraint on the $V$ and $A$ couplings is obtained:

$$|C_{e,V}C_{e,A}| \lesssim 10^{-8}, \quad (35)$$

which is automatically satisfied by our charge assignment.

Also, there are contributions from a $Z'$ to the anomalous magnetic moments of the electron and muon \[35–37\]. The one-loop contributions $\delta a_i$, mediated by a $Z'$, lead to

$$\delta a_e = 7.6 \times 10^{-6}C_{e,V}^2 - 3.8 \times 10^{-5}C_{e,A}^2 \quad (36)$$

$$-26 \times 10^{-13} \leq \delta a_e \leq 8 \times 10^{-13}, \quad (37)$$

$$|\delta a_\mu| = |0.009C_{\mu,V}^2 - C_{\mu,A}^2| \leq 1.6 \times 10^{-9}. \quad (38)$$

Another constraint is from electron-positron colliders (e.g., KLOE2) \[38\] through $e^+e^- \to \gamma Z'$, $Z' \to e^+e^-$. From this process one finds

$$\frac{C_{e,V}^2 + C_{e,A}^2}{\text{BR}(Z' \to e^+e^-)} \lesssim 3.7 \times 10^{-7}. \quad (39)$$

There is also a limit due to neutral pion decay, wherein the $V$ couplings of such a light state with quarks are, in general, strongly constrained from $\pi^0 \to Z' + \gamma$ searches at the NA48/2 experiment \[39\]. The process is proportional to the anomaly factor $N_\pi = \frac{1}{2}(2C_{u,V} + C_{d,V})^2$. Therefore, one gets the following bound:

$$\langle 0^+||\sigma_p||S\rangle = -0.132(33). \quad (31)$$
Finally, we discuss constraints arising from flavor-changing neutral currents (FCNCs). Despite an initially diagonal charge matrix, as the coupling strength between the first two generations and the third differs, rotations into the mass eigenstate will generate off-diagonal interactions, in the form of tree-level FCNCs. First, we examine \( K \to \pi e^+e^- \) via tree-level on-shell \( Z' \) exchange. Since we have a mass \( M_{Z'} \approx 17 \text{ MeV} \), one does not have contributions to \( K \to \pi \mu^+\mu^- \). There are strict limits here from LHCb [40]; however, there is no sensitivity to our \( Z' \) simply because the invariant mass range of \( e^+e^- \) begins from 20 MeV. This is done because the resolution degrades rapidly at small mass due to the background from photon conversion in the detector material. Future measurements may sample from smaller invariant masses, which could act as a discovery tool, or disprove our particular scenario. Next, we turn to \( B^0-B^{0*} \) mixing. As a first approximation, we use the results from Ref. [41], but assuming now a light \( Z' \), such that the propagator \( P \equiv (m_{Z'}^2 - M_{Z}^2)^{-1} \approx m_{Z}^2 \), rather than their approximation \( P \approx M_{Z}^{-2} \). This leads to the condition

\[
|g_{sb}^{(k)}| \lesssim 10^{-6},
\]

where [upon assuming minimal flavor violation in the quark sector and introducing Cabibbo-Kobayashi-Maskawa (CKM) matrix elements]

\[
g_{sb}^L = (V_{CKM}^{\dagger})_{13} (g_{Q_1}, g_{Q_3}, g_{Q_3}) V_{CKM} \approx 23, \tag{42}
\]

\[
g_{sb}^R = (V_{CKM}^{\dagger})_{13} (g_{u_R}, g_{u_R}, g_{u_R}) V_{CKM} \approx 23, \tag{43}
\]

and \( g_{i}^0 = g_{i} z_i \), for \( i = \{ Q_1, Q_3, u_R, d_R \} \). For our assignment, since it is family-universal in the \( L \)-handed quark sector, \( g_{sb}^L = 0 \), and only the \( K \)-handed sector, \( g_{sb}^R \propto V_{tb} V_{ts} \), contributes, this leads to the condition

\[
g', \tilde{g} \lesssim 10^{-4}. \tag{44}
\]

A similar estimate for the \( K-\bar{K} \) mixing yields a less stringent constraint. Despite a smaller propagator suppression (because \( m_K < m_B \)), the CKM suppression is now much stronger, \( \propto V_{td} V_{ts} \), and so one finds the weaker constraint \( g', \tilde{g} \lesssim 10^{-3} \).

In the scope of this paper, we do not perform a full flavor analysis of the \( B-\bar{B} \) and \( K-\bar{K} \) mixings, but we leave this as an approximate requirement.

One may expect constraints from the lepton sector also, such as from \( \tau \to 3\mu [42] \), but since the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix will depend on the particular nature of neutrino masses, which we do not specify here, we assume it to be always possible to construct the latter so as to avoid these kinds of lepton-flavor-violating limits.

\[
|2C_{u,V} + C_{d,V}| \lesssim \frac{0.3 \times 10^{-3}}{\sqrt{\text{BR}(Z' \to e^+e^-)}}. \tag{40}
\]

We now present the results for our particular charge assignment shown in Table I, consistent with all of the aforementioned experimental constraints.\(^4\) In Fig. 1, we plot the allowed parameters in the space of the \( U(1) \) gauge coupling, \( g' \), and the gauge-kinetic mixing strength, \( \tilde{g} \).

 FIG. 1. Allowed parameter space mapped on the \((g', \tilde{g})\) plane explaining the anomalous \(^8\text{Be} \) decay for \( Z' \) solutions with mass 16.7 (red), 17.3 (purple), and 17.6 (green) MeV. The white regions are excluded by the nonobservation of the same anomaly in the \(^8\text{Be} \) transition. Also shown are the constraints from \((g - 2)_\mu \) to be within the two dashed lines; \((g - 2)_e \), to be inside the two dotted lines (shaded in blue); and the electron beam dump experiment, NA64, to be in the shaded blue region outside the two solid lines. The surviving parameter space lies at small positive and negative \( \tilde{g} \) (though not at \( \tilde{g} = 0 \)), inside the dark shaded blue region which overlaps the Atomki anomaly solutions.

\(^4\) Other charge assignments are also possible by relaxing the conditions we impose.
The Atomki Collaboration solutions are also shown (in orange) in the NMEs of Eq. (31). One finds that the lower limit of \( Z \) in the Atomki anomaly for all three masses, shaded in red, purple, and green. Furthermore, the largest upper bound decreases with heavier \( a \) larger density of upper BR bounds at smaller values of it.\n
The combination of these two effects motivates why heavier \( M_{Z'} \) values have a larger range of solutions [i.e., a thicker green (17.6 MeV) than red (16.7 MeV) band in Fig. 1].

**IV. CONCLUSION**

In conclusion, with the assumption that the first two families of SM quark and lepton masses are generated by some high-scale physics, unlike those of the third family, which stem from a SM Higgs mechanism supplemented by an additional \( U(1)' \) (broken) group, yielding a very light \( Z' \) state, we have found a family-dependent (nonuniversal) charge assignment which can successfully accommodate the Atomki anomaly, in addition to all other experimental constraints on such a low-scale physics. This happens for a range of \( Z' \) masses (and corresponding decay rates), including the best fit of \( M_{Z'} = 16.7 \) MeV as well as other two published values, 17.3 and 17.6 MeV, over the coupling ranges \( g' \sim 10^{-5} \) and \( 1 \times 10^{-5} \lesssim |\tilde{g}| \lesssim 5 \times 10^{-5} \) for the gauge and kinetic mixing couplings, respectively, regulating the \( Z' \) interactions with SM fermions.

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