Dark matter and B -meson anomalies in a flavor dependent gauge symmetry

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A possibility of explaining the anomalies in the semileptonic B -meson decay $B \rightarrow K^* \mu \bar{\mu}$ has been explored in the framework of the gauged $U(1)_{\mu-\tau}$ symmetry. Apart from the muon anomalous magnetic moment and neutrino sector, we formulate the model starting with a valid Lagrangian and consider the constraints from the neutral meson mixings, the bounds on direct detection and the relic density of the bosonic dark matter candidate augmented to collider constraints. We search the parameter space, which accommodate the size of the anomaly of the $B \rightarrow K^* \mu \bar{\mu}$ decay, to satisfy all experimental constraints. We found the allowed region on the plane of the dark matter and Z' masses is rather narrow compared to the previous analysis.

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I. INTRODUCTION

A flavor dependent gauge symmetry is one of the promising candidates for new physics to describe the anomalies and other phenomenologies related with the flavor physics as well as to ensure the dark matter (DM) stability. In particular, the model of $U(1)_{\mu-\tau}$ provides several phenomenological prescriptions to resolve, namely, the muon anomalous magnetic moment [1],¹ experimental anomalies of semileptonic B -meson decay [2,3], the neutrino sector [4–15], and other related topics [16,17]. Among them, the B -meson decay anomaly is a very challenging topic due to some indications of new physics have been suggested in the B physics. For example, the angular observable P'_5 in the decay of the B meson $B \rightarrow K^* \mu^+ \mu^-$ [18] has been measured with a deviation of 3.4σ from the integrated luminosity of 3.0 fb^{-1} at the LHCb [19], which confirms the previous result with a deviation of 3.7σ [20]. In addition, the same observable was measured by Belle Collaboration [21,22] with the deviation

of 2.1σ . Furthermore, an anomaly in the measurement of the ratio of branching fraction $R_K = \text{BR}(B^+ \rightarrow K^+ \mu^+ \mu^-) / \text{BR}(B^+ \rightarrow K^+ e^+ e^-)$ [23,24] at the LHCb indicates a deviation of 2.6σ from the lepton universality predicted in the standard model (SM) [25]. Recently, the LHCb Collaboration has also measured the ratio of $R_{K^*} = \text{BR}(B \rightarrow K^* \mu^+ \mu^-) / \text{BR}(B \rightarrow K^* e^+ e^-)$, which deviates from the SM prediction by $\sim 2.4\sigma$ as $R_{K^*} = 0.660^{+0.110}_{-0.070} \pm 0.024(0.685^{+0.113}_{-0.069} \pm 0.047)$ for $(2m_\mu^2) < q^2 < 1.1 \text{ GeV}^2$ ($1.1 \text{ GeV}^2 < q^2 < 6 \text{ GeV}^2$) [26].

In a previous study of Ref. [2], we have proposed the flavor dependent gauge symmetry. This has successfully explained the anomaly of $B \rightarrow K^* \ell^+ \ell^-$ decay through generating the flavor violating Z' boson interactions at the one-loop level. In the model, the only Wilson coefficient C_9 with μ and τ in the B decays are generated via an extra Z' boson exchange, but this is explicitly not applicable for $B \rightarrow K^* e \bar{e}$ process [18]. We have also included a bosonic DM χ candidate and the vectorlike exotic quarks Q'_a , which are needed to generate the Wilson coefficient C_9 at the one-loop level. Thus, the DM relic density [27] can be explained by the measured anomalies in the decay of $B \rightarrow K^* \ell^+ \ell^-$ via the s-channel process mediated by the Z' boson exchange [28,29], where the Z' boson exchange can avoid a conflict with the constraints from the spin-independent DM direct detection searches, such as the experiments of LUX [30] and XENON1T [31].

In this paper, we adopt the flavor dependent gauge symmetry model with a more complete manner. We then perform a more detailed analysis in which we take into account the decay width of the Z' boson that is related to the relic density of the DM, the constraints from the

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¹Recently, a stringent constraint of the neutrino-trident process gives narrower parameter spaces of the extra gauge coupling (g') and mass ($m_{Z'}$).

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spin-independent DM-nucleon elastic scattering cross section mediated by the vectorlike quarks Q'_a at tree level, the large electron-positron (LEP) collider, and the large hadron collider (LHC), to improve the previous analyses of Ref. [2]. In our present numerical analysis, we find that the allowed regions of the DM and Z' masses are narrower than that of the previous analysis. This is expected because the decay width of Z' is larger.

This paper is organized as follows. In Sec. II, we briefly introduce a valid Lagrangian of our model, including the Higgs potential with the inert conditions, the B -meson anomaly, the collider physics, and the neutral meson mixings. A brief comment on how to directly produce the exotic quarks and the experimental constraints of the DM are also presented. In Sec. III, we present our numerical analysis results. Finally, Sec. IV is devoted to the summary of our results and conclusions.

II. MODEL SETUP AND CONSTRAINTS

In this section, we present a formulation of our model. We briefly introduce a gauged $U(1)_{\mu-\tau}$ symmetry with three families of the vectorlike isospin doublet quarks Q'_a , an isospin singlet inert complex boson χ , and singlet boson φ with a nonzero vacuum expectation value (VEV), which is denoted by $\langle\varphi\rangle \equiv v_\varphi/\sqrt{2}$, where H is the SM Higgs and its VEV is denoted by $\langle H\rangle \equiv v_H/\sqrt{2}$. The charge assignments of these new fermion and boson fields are summarized in Tables I and II, respectively.

TABLE I. Field contents of fermions and their charge assignments under $SU(2)_L \times U(1)_Y \times U(1)_{\mu-\tau}$, where $q_x \neq 0$ and the lower index a of Q' is the number of family which runs over 1–3.

Fermions	Leptons						Exotic vector fermions
	L_{L_e}	L_{L_μ}	L_{L_τ}	e_R	μ_R	τ_R	
$SU(3)_C$	1	1	1	1	1	1	3
$SU(2)_L$	2	2	2	1	1	1	2
$U(1)_Y$	$-\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	-1	-1	-1	$\frac{1}{6}$
$U(1)_{\mu-\tau}$	0	1	-1	0	1	-1	q_x

TABLE II. Field contents of bosons and their charge assignments under $SU(2)_L \times U(1)_Y \times U(1)_{\mu-\tau}$, where $q_{x,y} \neq 0$, $SU(3)_C$ singlet for all bosons, and χ is a complex boson that is considered as a DM candidate. Hence, $q_y \neq \pm q_x$, $q_y \neq \pm 2q_x$, and $q_x \neq \pm 3q_y$ are simultaneously satisfied.

Bosons	VEV $\neq 0$		Inert
	H	φ	
$SU(2)_L$	2	1	1
$U(1)_Y$	$\frac{1}{2}$	0	1
$U(1)_{\mu-\tau}$	0	q_y	q_x

A relevant Lagrangian under these symmetries is defined by

$$-\mathcal{L} = y_{\ell_i} \bar{L}_{L_i} H e_{R_i} + f_{aj} \bar{Q}'_{R_a} Q_{L_j} \chi + M_a \bar{Q}'_a Q'_a + \text{H.c.}, \quad (1)$$

where $a = 1-3$ and $(i, j) = e, \mu, \tau$ are generation indices, and the quark sector Q_{L_j} is the same as the SM. Note that the charged-lepton sector is diagonal due to the $U(1)_{\mu-\tau}$ symmetry.

The Higgs potential is given by

$$V = m_H^2 |H|^2 + m_\varphi^2 |\varphi|^2 + m_\chi^2 |\chi|^2 + \frac{1}{4} \lambda_H |H|^4 + \frac{1}{4} \lambda_\varphi |\varphi|^4 + \frac{1}{4} \lambda_\chi |\chi|^4 + \lambda_{H\varphi} |H|^2 |\varphi|^2 + \lambda_{H\chi} |H|^2 |\chi|^2 + \lambda_{\varphi\chi} |\varphi|^2 |\chi|^2, \quad (2)$$

where each field is defined to be

$$H = \begin{pmatrix} h^+ \\ \frac{1}{\sqrt{2}}(v_H + h + iz) \end{pmatrix}, \quad \varphi = \frac{1}{\sqrt{2}}(v_\varphi + \varphi_R + iz'), \quad (3)$$

where h^+ , z , and z' are respectively absorbed by the gauged bosons W^+ , Z , and Z' . After inserting the tadpole conditions for H and φ , the charge parity (CP)-even mass matrix is obtained as

$$M_{\text{even}}^2 = \begin{pmatrix} v_H^2 \lambda_H & 2v_H v_\varphi \lambda_{H\varphi} \\ 2v_H v_\varphi \lambda_{H\varphi} & v_\varphi^2 \lambda_\varphi \end{pmatrix}. \quad (4)$$

After the diagonalization, the mass eigenvalues and eigenstates are respectively given by

$$\begin{aligned} OM_{\text{even}}^2 O^T &= \begin{pmatrix} c_\theta & -s_\theta \\ s_\theta & c_\theta \end{pmatrix} \begin{pmatrix} v_H^2 \lambda_H & 2v_H v_\varphi \lambda_{H\varphi} \\ 2v_H v_\varphi \lambda_{H\varphi} & v_\varphi^2 \lambda_\varphi \end{pmatrix} \\ &\times \begin{pmatrix} c_\theta & s_\theta \\ -s_\theta & c_\theta \end{pmatrix} \\ &= \text{diag}(m_{h_{\text{SM}}}^2, m_H^2), m_{\text{hSM}}^2 \\ &= \frac{1}{2} \left(A - \sqrt{(v_\varphi^2 \lambda_\varphi - v_H^2 \lambda_H)^2 + 16v_\varphi^2 v_H^2 \lambda_{H\varphi}^2} \right), \\ m_H^2 &= \frac{1}{2} \left(A + \sqrt{(v_\varphi^2 \lambda_\varphi - v_H^2 \lambda_H)^2 + 16v_\varphi^2 v_H^2 \lambda_{H\varphi}^2} \right), \end{aligned} \quad (5)$$

where $A = v_\varphi^2 \lambda_\varphi + v_H^2 \lambda_H$, $\sin\theta(\cos\theta) \equiv c_\theta(s_\theta)$, and $\tan\theta \equiv t_\theta$ satisfy the following relation:

$$2v_H v_\varphi \lambda_{H\varphi} t_\theta^2 - (v_\varphi^2 \lambda_\varphi - v_H^2 \lambda_H) t_\theta - 2v_H v_\varphi \lambda_{H\varphi} = 0. \quad (6)$$

The mass eigenvalue of χ is given by

$$m_\chi^2 = v_\varphi^2 \lambda_{\varphi\chi} + v_H^2 \lambda_{H\chi} + 2m_\chi^2. \quad (7)$$

The inert conditions for χ are given by

$$\begin{aligned} \lambda_H, \lambda_\varphi, \lambda_\chi, \lambda_{H\varphi}, \lambda_{H\chi}, \lambda_{\varphi\chi} &> 0, \\ m_\chi^2 &> 0, \\ \lambda_H \lambda_\varphi &> 4\lambda_{H\varphi}^2. \end{aligned} \quad (8)$$

The Z' boson: We have a Z' boson from $U(1)_{\mu-\tau}$ gauge symmetry. After φ develops its VEV, the mass of the Z' boson is generated as

$$m_{Z'} = q_y^2 g^2 v_\varphi^2, \quad (9)$$

where g' is a gauge coupling for $U(1)_{\mu-\tau}$. The gauge interactions among Z' and fermions are given by

$$\begin{aligned} \mathcal{L} \supset g' Z'_\mu &\left(\bar{L}_{L_\mu} \gamma^\mu L_{L_\mu} - \bar{L}_{L_\tau} \gamma^\mu L_{L_\tau} + \bar{\mu}_R \gamma^\mu \mu_R - \bar{\tau}_R \gamma^\mu \tau_R \right. \\ &\left. + q_y \sum_{a=1}^3 \bar{Q}'_a \gamma^\mu Q'_a \right). \end{aligned} \quad (10)$$

Note that we ignore the kinetic mixing effects between $U(1)_{\mu-\tau}$ and $U(1)_Y$ by assuming that the contributions are relatively very small.

An explanation of the anomaly in $B \rightarrow K^{()} \ell^+ \ell^-$ decay:* In our case, we have a Wilson coefficient C_9 , which is associated with $(\bar{s} \gamma_\mu P_L b)(\mu \gamma^\mu \mu)$, via Fig. 1. Then, their contribution to a Wilson coefficient $\Delta C_9^{\mu\mu}$ is given by [2,18]

$$\begin{aligned} \Delta C_9^{\mu\mu} &\simeq \frac{q_x g^2}{(4\pi)^2 m_{Z'}^2 C_{\text{SM}}} \sum_{a=1}^3 f_{3a}^\dagger f_{a2} \int [dx] \ln \left(\frac{\Delta[M_a, m_\chi]}{\Delta[m_\chi, M_a]} \right), \\ C_{\text{SM}} &\equiv \frac{V_{tb} V_{ts}^* G_F \alpha_{\text{em}}}{\sqrt{2}\pi}, \end{aligned}$$

$$\Delta[m_1, m_2] = (x+y-1)(xm_b^2 + ym_s^2) + xm_1^2 + (y+z)m_2^2, \quad (11)$$

where $[dx] \equiv \int_0^1 dx dy dz \delta(1-x-y-z)$, V_{tb} , V_{ts} are the 3-3 and 3-2 elements of the Cabibbo-Kobayashi-Maskawa

FIG. 1. The diagrams introduce an effective coupling for the $Z'_\mu \bar{b} \gamma^\mu s + \text{H.c.}$ interaction.

(CKM) matrix, respectively, $G_F \approx 1.17 \times 10^{-5} [\text{GeV}]^{-2}$ is the Fermi constant, $\alpha_{\text{em}} \approx 1/137$ is the electromagnetic fine-structure constant, and $m_b \approx 4.18 \text{ GeV}$ and $m_s \approx 0.095 \text{ GeV}$ are the bottom and strange quark masses, respectively, which are given by the $\overline{\text{MS}}$ renormalization scheme at a scale $\mu = 2 \text{ GeV}$ [32]. We assume $m_b, m_s \ll m_{Z'}$ in Eq. (11). m_χ and M_a are respectively χ and Q'_a masses. We use the global fit for the value of C_9 [33], which gives the best fit value of $\Delta C_9 \approx -1.09$. The possible values of the ΔC_9 are respectively obtained as

$$\Delta C_9 = [-1.29, -0.87] \quad \text{for } 1\sigma, \quad [-1.67, -0.39] \quad \text{for } 3\sigma. \quad (12)$$

Constraint from K -meson decay: Here we discuss a constraint from K -meson decay. Considering a small effect of the charge parity (CP) violation emerges from new physics, the strongest bound is derived by $K_S \rightarrow \pi^0 \mu \bar{\mu}$ decay. It is then induced by the effective Hamiltonian [34,35]

$$H_{\text{eff}} = -\frac{G_F}{\sqrt{2}} V_{ud} V_{us}^* C_{7V}^{\mu\mu} (\bar{s} \gamma^\mu (1 - \gamma_5) d) (\bar{\mu} \gamma_\mu \mu). \quad (13)$$

Similar to the Wilson coefficient C_9 case, the Wilson coefficient ΔC_{7V} arising from the Z' exchange can be induced at the one-loop level. We then obtain

$$\begin{aligned} \Delta C_{7V}^{\mu\mu} &\simeq \frac{q_x g^2}{(4\pi)^2 m_{Z'}^2} \frac{1}{\sqrt{2} G_F V_{ud} V_{us}^*} \\ &\times \sum_{a=1}^3 f_{2a}^\dagger f_{a1} \int [dx] \ln \left(\frac{\Delta[M_a, m_\chi]}{\Delta[m_\chi, M_a]} \right). \end{aligned} \quad (14)$$

The values of the experimental constraint are given by [34]

$$|a_S^{\mu\mu}| = 1.54_{-0.32}^{+0.40}, \quad a_S^{\mu\mu} \equiv \frac{2\pi}{\alpha_{\text{em}}} V_{ud} V_{us}^* C_{7V}^{\mu\mu}, \quad (15)$$

where a large uncertainty is found. For estimating the Z' contribution, we derive

$$\begin{aligned} \Delta a_S^{\mu\mu} &\simeq 0.46 \left(\frac{g'}{0.1} \right)^2 \left(\frac{100 \text{ GeV}}{m_{Z'}} \right)^2 \\ &\times \sum_{a=1}^3 f_{2a}^\dagger f_{a1} \int [dx] \ln \left(\frac{\Delta[M_a, m_\chi]}{\Delta[m_\chi, M_a]} \right). \end{aligned} \quad (16)$$

We find that $\Delta a_S^{\mu\mu}$ lies sufficiently within the range of the experimental uncertainty if the Yukawa coupling f is taken $O(0.1)$ and the constraint from Δm_K tends to be stronger.

The LHC constraint: Since we consider no interactions between the electron-positron pair and Z' , we evade the stringent constraint from the LEP experiment. However, we need a constraint from the LHC experiment, while we

consider interactions among u and d quarks via one-loop contributions, as it can be explicitly seen in the part of $b \rightarrow s\ell\bar{\ell}$.

The most stringent constraint from the LHC arises from the process of $q\bar{q} \rightarrow Z' \rightarrow \bar{\mu}\mu$. From this constraint, the effective mass bound suggests $30 \text{ TeV} \lesssim \Lambda$, where its effective operator is given by $\frac{1}{\Lambda^2}(\bar{q}\gamma_\mu P_L q)(\bar{\mu}\gamma^\mu P_L \mu)$ [36]. We then can easily estimate the bound on $m_{Z'}$ in the effective operator analysis. Similar to the case of $b \rightarrow s\bar{\mu}\mu$, our effective operator is defined as

$$\mathcal{L} \simeq \frac{q_x g^2}{(4\pi)^2 m_{Z'}^2} \sum_{a=1}^3 f_{1a}^\dagger f_{a1} \int [dx] \ln \left(\frac{xM_a^2 + (y+z)m_\chi^2}{xm_\chi^2 + (y+z)M_a^2} \right) \times (\bar{q}\gamma_\mu P_L q)(\bar{\mu}\gamma^\mu P_L \mu), \quad (17)$$

where $q = (u, d)$. It implies that we have the following constraint from the LHC:

$$\frac{m_{Z'}}{g'} \gtrsim 30 \left[\frac{q_x}{(4\pi)^3} \sum_{a=1}^3 f_{1a}^\dagger f_{a1} \int [dx] \ln \left(\frac{xM_a^2 + (y+z)m_\chi^2}{xm_\chi^2 + (y+z)M_a^2} \right) \right]^{\frac{1}{2}} \times \text{TeV}. \quad (18)$$

When we take a degenerate mass for M_a as 1000 GeV and $m_\chi = 100$ GeV, the constraint then gives

$$\frac{m_{Z'}}{g'} \gtrsim 0.45 \left[\sum_{a=1}^3 f_{1a}^\dagger f_{a1} \right]^{\frac{1}{2}} \text{TeV}. \quad (19)$$

This shows that the constraint can be easily avoided while the coupling f_{1a} is not too large.

M – $\bar{M}$ mixing: The neutral meson mixings also give the constraints of the parameter space. The neutral meson mixings are shown in Fig. 2 and their formulas are lower than the experimental bounds as follows [32,37]:

$$\Delta m_K \approx \sum_{a,b=1}^3 f_{2a}^\dagger f_{a1} f_{2b}^\dagger f_{b1} F_{box}^K[M_a, M_b] \lesssim 3.48 \times 10^{-15} \text{ [GeV]}, \quad (20)$$

FIG. 2. The diagram representing the neutral meson mixings.

$$\Delta m_{B_d} \approx \sum_{a,b=1}^3 f_{3a}^\dagger f_{a1} f_{3b}^\dagger f_{b1} F_{box}^{B_d}[M_a, M_b] \lesssim 3.36 \times 10^{-13} \text{ [GeV]}, \quad (21)$$

$$\Delta m_{B_s} \approx \sum_{a,b=1}^3 f_{3a}^\dagger f_{a2} f_{3b}^\dagger f_{b2} F_{box}^{B_s}[M_a, M_b] \lesssim 1.17 \times 10^{-11} \text{ [GeV]}, \quad (22)$$

$$\Delta m_D \approx \sum_{a,b=1}^3 f_{1a}^\dagger f_{a2} f_{1b}^\dagger f_{b2} F_{box}^D[M_a, M_b] \lesssim 6.25 \times 10^{-15} \text{ [GeV]}, \quad (23)$$

$$F_{box}^M(m_1, m_2) = \frac{m_M f_M^2}{3(4\pi)^2} \int_0^1 \frac{x[dx]}{xm_\chi^2 + ym_1^2 + zm_2^2}, \quad (24)$$

where m_M and f_M are the meson mass and the meson decay constant, respectively. The following parameter values are used in our analysis: $f_K \approx 0.156$ GeV, $f_{B_d(B_s)} \approx 0.191(0.274)$ GeV [38,39],² $f_D \approx 0.212$ GeV, $m_K \approx 0.498$ GeV, $m_{B_d(B_s)} \approx 5.280(5.367)$ GeV, and $m_D \approx 1.865$ GeV.

Constraints from direct production of Q' 's: The exotic quarks Q' 's can be produced in a pair via QCD processes at the LHC. Each Q' will then decay into $Q' \rightarrow q_i \chi$, where q_i represents a quark with flavor i . Hence, searching for “ $\{tt, bb, tj, bj, jj\} + \text{missing } E_T$ ” signals will constrain our present model. The branching ratios for a particular quark flavor i depend on the relative sizes of the Yukawa couplings, f_{3j} and f_{aj} with $a = 1, 2$. We then roughly estimate the lower limit on the mass of Q' using the current LHC data for the s -quark searches [40,41], which indicates the mass should be larger than $\sim 0.5\text{--}1$ TeV, which depends on the mass difference between Q' and χ . In our following analysis, we simply take the value of $M_a > 1$ TeV in order to satisfy this constraint.

Dark matter: In our model scenario, a complex scalar χ is considered as a DM candidate, since we have a remnant Z_2 symmetry after the $\mu - \tau$ symmetry breaking. The DM candidate and Q'_a are odd under the Z_2 symmetry and the other particles are even. If the M_a is heavier than the m_χ , the DM candidate is a stable particle. The DM dominantly annihilates into the SM leptons via $\chi\chi \rightarrow Z' \rightarrow \mu^+ \mu^- (\tau^+ \tau^-)$,³ so that the DM in our model is naturally leptophilic. The relic density of the DM is given by

²We thank Alexander Lenz for bringing up a new value of f_{B_s} , which includes a bag parameter dependence.

³Notice that we do not rely on the Higgs portal, although there are two resonant solutions at around the half masses of the SM Higgs and another neutral Higgs [42].

FIG. 3. A Feynman diagram contributing to the spin-independent DM-nucleon elastic scattering.

$$\Omega h^2 \approx \frac{1.07 \times 10^9}{\sqrt{g_*(x_f)} M_{\text{Pl}} J(x_f) [\text{GeV}]}, \quad (25)$$

where $g^*(x_f \approx 25) \approx 100$, $M_{\text{Pl}} \approx 1.22 \times 10^{19}$, and $J(x_f)$ ($\equiv \int_{x_f}^{\infty} dx \frac{(\sigma v_{\text{rel}})}{x^2}$) is given by

$$J(x_f) = \int_{x_f}^{\infty} dx \left[\frac{\int_{4m_\chi^2}^{\infty} ds \sqrt{s - 4m_\chi^2} (\sigma v_{\text{rel}}) K_1\left(\frac{\sqrt{s}}{m_\chi} x\right)}{16m_\chi^5 x [K_2(x)]^2} \right],$$

$$(\sigma v_{\text{rel}}) = \frac{g'^4 q_x^2 s (s - m_\chi^2)}{3\pi [(s - m_{Z'}^2)^2 + m_{Z'}^2 \Gamma_{Z'}^2]},$$

$$\Gamma_{Z'} \approx \frac{g'^2 m_{Z'}}{4\pi} + \frac{|g' q_x|^2}{16\pi m_{Z'}^2} (m_{Z'}^2 - 4m_\chi^2)^{3/2}. \quad (26)$$

With s as a Mandelstam variable, and $K_{1,2}$ are the modified Bessel functions of the second kind of order 1 and 2, respectively, we expect Z' decays into $\mu\bar{\mu}$, $\tau\bar{\tau}$, and $\chi\chi^*$ pairs.⁴ In our numerical analysis, we use the current experimental range for the relic density at 3σ confidential level [27]: $0.11 \lesssim \Omega h^2 \lesssim 0.13$.

Direct detection of DM: The dominant elastic scattering cross section arises from the Q' exchange process in Fig. 3, and its effective Lagrangian of the component level is given by

$$\mathcal{L} \simeq -i \sum_{i=1}^3 \sum_{a=1}^3 \frac{f_{ia} f_{ai}}{2M_a^2} [\bar{q}_i \gamma^\mu P_L q_i] [\chi^* \overleftrightarrow{\partial}_\mu \chi], \quad (27)$$

where we use the following assumptions of the four transferred momentum $q^2 = (p_1 - p_2)^2 \ll M_a^2$ and the nucleus of a target almost stops (at rest frame); $\chi^*(p_1) (p_1 - p_2)_\mu \chi(k_1) = \frac{1}{2} \chi^*(p_1) [(p_1 - p_2) + (k_1 - k_2)]_\mu \chi(k_1) \approx \frac{1}{2} \chi^*(p_1) (p_1 + k_1)_\mu \chi(k_1) = \frac{i}{2} \chi^* \overleftrightarrow{\partial}_\mu \chi$, where the right and left

⁴Without decay width of Z' , the cross section at around the pole of $2m_\chi = m_{Z'}$ is too large and the relic density would be underestimated.

sides correspond to the operators in momentum space and spacetime, respectively. We then straightforwardly define the DM-nucleon elastic scattering operator as follows:

$$\mathcal{L}_N \simeq -i \sum_{i=u}^{d,s} \sum_{a=1}^3 \frac{f_{ia} f_{ai}}{4M_a^2} [F_1^{q_i/N} \bar{N} \gamma^\mu N - F_A^{q_i/N} \bar{N} \gamma^\mu \gamma_5 N] \times (\chi^* \overleftrightarrow{\partial}_\mu \chi), \quad (28)$$

where we assume that this process is an elastic scattering, then the four transferred momentum is expressed by $q^2 \approx 0$, where p_1 and p_2 are the four momentum of the χ^* and q_i or N , respectively, and $F_{1,A}^{q_i/N}$ are the form factors, which are taken from Ref. [43]. The squared matrix element is given by

$$|\mathcal{M}|^2 = \frac{1}{16} \left| \sum_{i=1}^3 \sum_{a=1}^3 \frac{f_{ia} f_{ai}}{M_a^2} (F_1^{q_i/N} \langle N(k_2) | \bar{N} \gamma^\mu N | N(p_2) \rangle - F_A^{q_i/N} \langle N(k_2) | \bar{N} \gamma^\mu \gamma_5 N | N(p_2) \rangle) \right|^2$$

$$\times |\langle \chi(p_1) | \chi^* \overleftrightarrow{\partial}_\mu \chi | \chi(k_1) \rangle|^2$$

$$\simeq 2m_N^2 m_\chi^2 \sum_{i=1}^3 \sum_{a=1}^3 \left(\left| \frac{f_{ia} f_{ai} F_1^{q_i/N}}{M_a^2} \right|^2 + \left| \frac{f_{ia} f_{ai} F_A^{q_i/N}}{M_a^2} \right|^2 \right), \quad (29)$$

where the first and second terms in Eq. (29) do not have an interference term.

Finally, the complete form of the DM-nucleon elastic scattering cross section is expressed by

$$\sigma \approx \left(\frac{m_\chi m_N}{m_N + m_\chi} \right)^2 \frac{|\mathcal{M}|^2}{32\pi m_\chi^2 m_N^2}$$

$$= \frac{1}{16\pi} \left(\frac{m_\chi m_N}{m_N + m_\chi} \right)^2$$

$$\times \sum_{i=1}^3 \sum_{a=1}^3 \left(\left| \frac{f_{ia} f_{ai} F_1^{q_i/N}}{M_a^2} \right|^2 + \left| \frac{f_{ia} f_{ai} F_A^{q_i/N}}{M_a^2} \right|^2 \right), \quad (30)$$

where $\sum F_1 \approx 3$ corresponds to the effective operator $(\bar{N} \gamma^\mu N) (\chi^* \overleftrightarrow{\partial}_\mu \chi)$, $\sum F_A \approx 0.49$ corresponds to the effective operator $(\bar{N} \gamma^\mu \gamma_5 N) (\chi^* \overleftrightarrow{\partial}_\mu \chi)$, and $m_N \approx 0.939$ GeV. The current experimental upper bounds for the cross section of the spin-independent DM-nucleon elastic scattering are, respectively, $\sigma_{\text{exp}} \lesssim 2.2 \times 10^{-46}$ cm² at $m_\chi \approx 50$ GeV for the LUX data [30] and $\sigma_{\text{exp}} \lesssim 4.1 \times 10^{-47}$ cm² at $m_\chi \approx 30$ GeV for the XENON data [31]. In our numerical analysis, we conservatively restrict the LUX/XENON1T bounds for the whole range of the DM mass.

III. NUMERICAL ANALYSIS

In this section, results for our numerical analysis are presented. In our analysis, we fix a parameter $|q_x| = 1$ for simplicity. The ranges of the input parameters are set as follows:

$$g' \in [0.1, 1], \quad f \in [0.001, \sqrt{4\pi}], \quad m_{Z'} \in [10, 300] \text{ [GeV]}, \\ m_\chi \in [10, 150] \text{ [GeV]}, \quad M_a \in [1000, 3000] \text{ [GeV]}, \quad (31)$$

where the lowest DM mass of 10 GeV is an assumption to satisfy a condition $(m_\tau/m_\chi)^2 \ll 1$ in the cross section of relic density. In this calculation, we assume $M_1 < M_2 < M_3$, and $m_{Z'} > m_\chi$. We then search the allowed region using the range of the input parameters listed above to satisfy all constraints, namely, $M - \bar{M}$ mixing, the measure of the relic density of the DM, the spin-independent DM-nucleon scattering cross section via the Z' boson exchange,

FIG. 4. The allowed region of ΔC_9 on the plane of m_χ and $m_{Z'}$ satisfy the relic density $0.11 \sim 0.13$ and the neutral meson mixings, $b \rightarrow s\gamma$. The blue, red, and green dots are, respectively, the region of $\Delta C_9 \lesssim -5$, $-1.29 \lesssim \Delta C_9 \lesssim -0.87$ at 1σ , and $-1.67 \lesssim \Delta C_9 \lesssim -0.39$ at 3σ .

FIG. 5. The allowed region on the $\Delta m_K - \Delta m_{B_s}$ plane (the left panel) and $\Delta m_{B_d} - \Delta m_{B_D}$ plane (the right panel), where the color representations are similar as in Fig. 4.

and the constraint of the LHC as well as explaining the anomaly of $b \rightarrow s\bar{\mu}\mu$ decay.

Figure 4 shows the allowed region on the plane of m_χ and $m_{Z'}$, where the blue, red, and green dots represent, respectively, the region corresponding to no constraint on ΔC_9 , 1σ range $-1.29 \lesssim \Delta C_9 \lesssim -0.87$, and 3σ range $-1.67 \lesssim \Delta C_9 \lesssim -0.39$. Note that the lowest point of the maximum absolute value of f is of the order 0.1 for $1(3)\sigma$. The correlation between m_χ and $m_{Z'}$ in Fig. 4 comes from the closed resonant point of the relic density of the DM when m_χ is heavier. In the lighter region of m_χ , the allowed region becomes wider due to the larger cross section. At around the resonant region of $10 \text{ GeV} \lesssim m_\chi \lesssim 40 \text{ GeV}$, there are no allowed regions since the corresponding cross section is too large to satisfy the relic density. This clearly indicates that the mass ranges of the DM and Z' are, respectively, $10 \text{ GeV} \leq m_\chi \lesssim 146 \text{ GeV}$ and $10 \text{ GeV} \leq m_{Z'} \lesssim 295 \text{ GeV}$, where the specific ranges mainly originates from the constraint of the relic density, although the lowest bound of DM mass of 10 GeV comes from the lowest input parameter of the DM mass. The anomaly of the decay of $b \rightarrow s\bar{\mu}\mu$ is well explained for the whole allowed region on the plane of m_χ and $m_{Z'}$.

In Fig. 5, we clearly show that the allowed regions on the planes of $\Delta m_K - \Delta m_{B_s}$ (the left panel) and $\Delta m_{B_d} - \Delta m_{B_D}$ (the right panel), where the color representation is similar as in Fig. 4. Figure 5 indicates the allowed region for explaining the anomaly of $b \rightarrow s\bar{\mu}\mu$, which tends to lie in the range of the experimental bounds in our parameter space. A branching ratio of $\text{BR}(b \rightarrow s\gamma)$ is restricted to be 4.02×10^{-4} , but the typical value is at most of the order 10^{-10} . Therefore, the present model clearly satisfies this constraint. Additionally, we found that the LHC constraint tends to be weaker than the experimental bounds for the neutral meson mixings in our parameter space.

Figure 6 shows the cross section of the nucleon-DM elastic scattering obtained from our parameter space scanned. This indicates the parameter points excluded by

FIG. 6. The nucleon-DM elastic scattering generated by the parameter space scanned. The current upper bound by XENON1T is represented by the solid line.

the present XENON1T data [31]. This means more parameter points can be explored in future experiments.

IV. SUMMARY AND CONCLUSIONS

We have explored the possibility of explaining the experimental anomalies in the semileptonic decay of the B -meson, $B \rightarrow K^* \mu \bar{\mu}$, in the framework of the gauged $U(1)_{\mu-\tau}$ symmetry. With our present model, which is built in a more complete manner than previous model, we have performed a more detailed analysis by searching the allowed region from several present experimental constraints.

Apart from the muon anomalous magnetic moment and neutrino sector, we have formulated a model starting with a valid Lagrangian by considering the Higgs potential with the inert conditions, the Wilson coefficient for the decay of $B \rightarrow K^* \mu \bar{\mu}$, the collider physics, the neutral meson

mixings, the bound on direct detection, and the relic density of a bosonic dark matter candidate. We have searched the parameter space, which explain the size of the anomaly of $B \rightarrow K^* \mu \bar{\mu}$ decay, satisfying all constraints. We found that the allowed region on the plane of the DM and Z' masses is narrower compared to the previous analysis in the heavier DM mass. This is expected because the decay width of Z' is rather large. On the other hand, in the lighter region of m_χ , the allowed region becomes wider because the cross section is larger. Moreover, there are no allowed region at around the resonant region of $10 \text{ GeV} \lesssim m_\chi \lesssim 40 \text{ GeV}$, since the corresponding cross section is too large to satisfy the relic density. The meson mixing of ΔB_s can be well tested in future experiments, since the structure of the Yukawa couplings f is the same as the one of ΔC_9 of $B \rightarrow K^* \mu \bar{\mu}$. The absolute parameter of f can naturally be estimated to explain ΔC_9 , and its minimum absolute is at most of the order ~ 0.1 .

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