Rare radiative $\Xi_b^- \to \Xi^- \gamma$ decay in the relativistic quark model

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Recently the LHCb collaboration put the upper limit on the $\Xi_b^- \to \Xi^- \gamma$ decay branching ratio $Br(\Xi_b^- \to \Xi^- \gamma) < 1.3 \times 10^{-4}$. The measured value is below the light-cone sum rule prediction. In this paper the rare radiative decay of the $\Xi_b^-$ baryon is studied in the framework of the relativistic quark model based on the quasipotential approach and QCD. The decay form factors are calculated with the comprehensive account of the relativistic effects. The obtained result for the branching ratio is found to be below the upper limit set by LHCb and is consistent with theoretical predictions based on the SU(3) flavor-symmetry, light-front quark model and light-cone QCD sum rules in full theory within theoretical uncertainties.

I. INTRODUCTION

In the standard model (SM) the exclusive rare weak decays of hadrons governed by the $b \to s$ quark transitions proceed through the flavor-changing neutral currents (FCNC). Thus such processes are forbidden at the tree level in SM. The leading contribution comes from the one-loop, so-called penguin diagrams. As a result, such a decay channel is strongly suppressed, which complicates its experimental search. The rare radiative decay $\Xi_b^- (bsd) \to \Xi^- (ssd) \gamma$ has not yet been observed experimentally. However, in 2021 the LHCb collaboration at Large Hadron Collider (CERN) set an upper experimental limit on its branching ratio [1]. Note that the branching ratio of the similar rare radiative decay $\Lambda_b (bud) \to \Lambda (sud) \gamma$ was measured by the LHCb Collaboration in 2019 [2].

Theoretically the rare radiative decay of the $\Xi_b^-$ baryon has been studied with different approaches. The prediction based on the $SU(3)$-flavor symmetry, which relates the decay branching fraction of the $\Xi_b$ baryon to the measured branching fraction of the $\Lambda_b$ baryon, gives the value consistent with the upper experimental limit [3]. The light-front quark model [4] and light-cone QCD sum rules within full theory predict the value satisfying this upper limit too [5], while the computation using light-cone sum rules shows a significant tension with the experimental value [6]. Therefore a more detailed theoretical investigation of this decay is required.

In this paper we comprehensively investigate the rare radiative decays $\Xi_b^- \to \Xi^- \gamma$ in the framework of the relativistic quark-diquark model based on the quasipotential approach. All relativistic effects, including the wave function transformations from the rest to the moving reference frame and contributions of the intermediate negative-energy states, are systematically taken into account. Using baryon wave functions, found in the previous studies of the baryon spectroscopy, we calculate the form factors parameterizing the baryon decay matrix element. The obtained form factors of the $\Xi_b^-$ baryon transitions are used for the prediction of the $\Xi_b^- \to \Xi^- \gamma$ decay branching ratio. Our result is consistent with the values from Refs. [3][5] within theoretical uncertainties and shows a small deviation in central values. Note that all these predictions are lower than experimental limit.
There are good chances that this decay will be soon observed by the LHCb Collaboration. So it will be possible to compare results obtained with different methods and decide which method gives a more precise prediction.

II. EFFECTIVE HAMILTONIAN

The effective Hamiltonian for the rare $b \rightarrow s$ transitions is given by [7]:

$$H_{\text{eff}} = -\frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \sum_{i=1}^{6} C_i(\mu) Q_i(\mu) + C_{7\gamma}(\mu) Q_{7\gamma}(\mu) + C_{SG}(\mu) Q_{SG}(\mu)$$

$$\equiv -\frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* Q^{T}(\mu) C(\mu)$$ \hspace{1cm} (1)

Here $C_i(\mu)$ are the Wilson coefficients, $G_F$ is the Fermi coupling constant, $V_{tb}$ and $V_{ts}$ are the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements, $Q_i(\mu)$ are the local operators given by

$$Q_1 = (\bar{s}_\alpha c_\beta)_{V-A} (\bar{c}_\beta b_\alpha)_{V-A},$$

$$Q_2 = (\bar{s}_\alpha c_\beta)_{V-A} (\bar{c}_\beta b_\beta)_{V-A},$$

$$Q_3 = (\bar{s}_\alpha b_\beta)_{V-A} \sum_q (\bar{q}_\beta q_\beta)_{V-A},$$

$$Q_4 = (\bar{s}_\beta b_\alpha)_{V-A} \sum_q (\bar{q}_\alpha q_\beta)_{V-A},$$

$$Q_5 = (\bar{s}_\alpha b_\alpha)_{V-A} \sum_q (\bar{q}_\alpha q_\beta)_{V+A},$$

$$Q_6 = (\bar{s}_\beta b_\alpha)_{V-A} \sum_q (\bar{q}_\alpha q_\beta)_{V+A},$$

$$Q_{7\gamma} = \frac{e}{4\pi^2} \bar{s}_\alpha \sigma^{\mu\nu} (m_b R + m_s L) b_\alpha F_{\mu\nu},$$

$$Q_{SG} = \frac{g_s}{4\pi^2} \bar{s}_\alpha \sigma^{\mu\nu} (m_b R + m_s L) T_{a\alpha\beta b\beta}^a C_{\mu\nu},$$

where $(\bar{q}_\alpha q_\beta)_{V\pm A} = \bar{q}_\alpha \gamma_\mu (1 \pm \gamma_5) q_\beta$, $Q^{T} = (Q_1, Q_2, ..., Q_{SG})$, $C^{T} = (C_1, C_2, ..., C_{SG})$, $\alpha$ and $\beta$ are color indices, $R = \frac{1+\gamma_5}{2}$ and $L = \frac{1-\gamma_5}{2}$, $\gamma_\mu$ and $\gamma_5$ are the Dirac matrices, $e$ and $g$ are electromagnetic and strong coupling constants, respectively. $F_{\mu\nu}$ is the electromagnetic field strength tensor, which in the case of the plane electromagnetic wave is given by

$$F_{\mu\nu} = -i(\epsilon_{\mu} q_{\nu} - \epsilon_{\nu} q_{\mu}) e^{i q x},$$

where $\epsilon_{\mu}$ is the polarization four-vector, $q_{\mu}$ is the four-momentum vector, and $\sigma^{\mu\nu}$ is a commutator of the Dirac matrices

$$\sigma^{\mu\nu} = \frac{i}{2} [\gamma_\mu \gamma_\nu - \gamma_\nu \gamma_\mu].$$

For the $b \rightarrow s \gamma$ transition the main contribution comes from the magnetic penguin operator $Q_{7\gamma}$, thus the effective Hamiltonian takes the form
$$H^{\text{eff}} = -\frac{G_F e^2}{4\pi^2 \sqrt{2}} V_{tb} V_{ts}^* C_{7\gamma}^{\text{eff}} (m_b) \bar{s} \sigma_{\mu \nu} \left[ m_b R + m_s L \right] b F^{\mu \nu}. \quad (2)$$

We first calculate the value of the relevant effective Wilson coefficient $C_{7\gamma}^{\text{eff}}$ in the leading order. To achieve this goal we need to solve the system of renormalization group equations

$$\frac{d C^{(0)\text{eff}}(\mu)}{d \ln \mu} = \frac{\alpha_s}{4\pi} (\hat{\gamma}^{(0)\text{eff}}) T C^{(0)\text{eff}}(\mu), \quad (3)$$

where index (0) stands for leading order, $\alpha_s(\mu)$ is the running strong coupling constant

$$\alpha_s(\mu) = \frac{\alpha_s(M_Z)}{1 - \frac{\beta_0}{2\pi} \alpha_s(M_Z) \ln \frac{M_Z}{\mu}}, \quad (4)$$

$$\beta_0 = \frac{11 N_c - 2 n_f}{3} \quad (5)$$

with the number of colors $N_c = 3$, and the number of quark flavors $n_f = 5$. We take the current world averaged value [12] of

$$\alpha_s(M_Z) = 0.1179 \pm 0.0010. \quad (6)$$

The effective anomalous dimension matrix $\hat{\gamma}^{(0)\text{eff}}$ is given by [7], [8]

$$\hat{\gamma}^{(0)\text{eff}} = \begin{pmatrix}
-2 & 6 & 0 & 0 & 0 & 0 & 0 & 0 & 3 \\
6 & -2 & -2 & -2 & -2 & -2 & 2 & 416 & 70 \\
0 & 0 & -22 & 22 & -4 & -4 & 4 & -464 & 545 \\
0 & 0 & 44 & 4 & -10 & -10 & 10 & 136 & 512 \\
0 & 0 & 0 & 0 & 2 & -6 & 32 & -59 & 9 \\
0 & 0 & 10 & 10 & -10 & -10 & 38 & -296 & -703 \\
0 & 0 & 0 & 0 & 0 & -32 & 9 & 28 & 3 \\
0 & 0 & 0 & 0 & 0 & 0 & -32 & 9 & 28 & 3
\end{pmatrix} \quad (7)$$

The initial conditions are as follows [7]

$$\begin{align*}
C_1^0(M_W) &= C_3^0(M_W) = C_4^0(M_W) = C_5^0(M_W) = C_6^0(M_W) = 0, \\
C_2^0(M_W) &= 1, \\
C_{7\gamma}^{(0)\text{eff}}(M_W) &= \frac{3 x^2 - 2 x^2}{4(x-1)^3} \ln x + \frac{-8 x^3 - 5 x^2 + 7 x}{24(x-1)^5} \approx -0.194, \\
C_{8G}^{(0)\text{eff}}(M_W) &= -\frac{3 y^2}{4(x-1)^4} \ln x + \frac{-x^3 + 5 x^2 + 2 x}{8(x-1)^3} \approx -0.097, 
\end{align*} \quad (8)$$

where $x \equiv \frac{m_t^2}{M_W^2} \approx 4.62$. 
Solving Eq. (3) with the initial conditions (8) we obtain the expression for the effective Wilson coefficient $C_{7\gamma}^{(0)\text{eff}}$

$$C_{7\gamma}^{(0)\text{eff}}(\mu) = \frac{16}{3} C_{7\gamma}^{(0)}(M_W) + \frac{8}{3} \left( \eta_{i}^{H} - \eta_{i}^{\omega} \right) C_{8G}^{(0)}(M_W) + C_{2}^{(0)}(M_W) \sum_{j=1}^{8} K_j \eta_j \gamma, \quad (9)$$

where

$$K_j = \left( -\frac{3}{7}, -\frac{1}{14}, -0.6494, -0.0380, 0.0185, -0.0057, 2.2996, -1.0880 \right), \quad (10)$$

$$c_j = \left( \frac{6}{23}, -\frac{12}{23}, 0.4086, -0.4230, -0.8994, 0.1456, \frac{14}{23}, \frac{16}{23} \right), \quad (11)$$

and

$$\eta \equiv \frac{\alpha_s(\mu m_W)}{\alpha_s(\mu)}.$$

Substituting numerical values we get the following result for the effective Wilson coefficient $C_{7\gamma}^{(0)\text{eff}}$ at $\mu = m_b$

$$C_{7\gamma}^{(0)\text{eff}}(m_b) = 0.674 C_{7\gamma}^{(0)}(M_W) + 0.091 C_{8G}^{(0)}(M_W) - 0.170 C_{2}^{(0)}(M_W) = -0.310. \quad (12)$$

### III. RELATIVISTIC QUARK MODEL

Now we calculate the matrix element of the effective Hamiltonian $\mathcal{H}^{\text{eff}}$ between the initial and final states

$$M = \langle \Xi^- | \mathcal{H}^{\text{eff}} | \Xi_b^- \rangle. \quad (13)$$

Note that in the absence of the QCD corrections we can make the following replacement in $\mathcal{H}^{\text{eff}}$: $-\sigma^{\mu\nu} F_{\mu\nu} \to 2i \sigma^{\mu\nu} \epsilon_{\mu} q_{\nu}$. As a result, the operator $Q_{7\gamma}$ reduces to $-2i \frac{\eta}{f^{\gamma}} s_{\gamma} \sigma^{\mu\nu} \epsilon_{\nu}(m_b R + m_s L) b_{\gamma} \gamma$. Thus to find $M$ we need to calculate the following matrix elements between baryon states

$$\langle \Xi^- (P) | \bar{s} \sigma^{\mu\nu} q_{\nu} b | \Xi_b^- (Q) \rangle, \quad (14)$$

and

$$\langle \Xi^- (P) | \bar{s} \sigma^{\mu\nu} q_{\nu} \gamma_5 b | \Xi_b^- (Q) \rangle. \quad (15)$$

The matrix elements (14) and (15) can be parameterized by the following set of form factors [9]

$$\langle \Xi^- | \bar{s} \sigma^{\mu\nu} q_{\nu} b | \Xi_b^- \rangle = \bar{u}_{\Xi^-}(p_{\Xi}, s) \left[ \frac{f^{\gamma}_T(q^2)}{m_{\Xi_b}} \left( \gamma^\mu q^\nu - \delta^\mu q^\nu \right) - f_2^{\gamma}(q^2) \right] \left( \gamma^\mu q^\nu - \delta^\mu q^\nu \right) u_{\Xi^-}(p_{\Xi_b}, s'), \quad (16)$$

$$\langle \Xi^- | \bar{s} \sigma^{\mu\nu} q_{\nu} \gamma_5 b | \Xi_b^- \rangle = \bar{u}_{\Xi^-}(p_{\Xi}, s) \left[ \frac{g^{\gamma}_T(q^2)}{m_{\Xi_b}} \left( \gamma^\mu q^\nu - \delta^\mu q^\nu \right) - g_2^{\gamma}(q^2) \right] \gamma_5 u_{\Xi^-}(p_{\Xi_b}, s'). \quad (17)$$

Since we investigate the decay with the emission of the real photon, we only need the values of the form factors $f_2^{\gamma}(q^2)$ and $g_2^{\gamma}(q^2)$ at $q^2 = 0$. 
Summing the expressions (16) and (17) we get the following parameterization for the matrix element of the effective Hamiltonian between baryon states

\[
\langle \Xi^{-}\gamma|H_{\text{eff}}|\Xi^{-}_{b}\rangle = \frac{G_F m_b e}{4\pi^2 \sqrt{2}} V_{tb} V_{ts} C_{\gamma}^{(0)\text{eff}} (m_b) \bar{u}_{\Xi^{-}}(p_{\Xi^{-}}, s) i\sigma^\mu \epsilon_\nu q_\nu \left( g_\nu f_2^T(0) + \gamma_5 g_A g_2^T(0) \right) u_{\Xi^{-}_b}(p_{\Xi^{-}_b}, s^\prime),
\]

where \( g_\nu = 1 + \frac{m_\pi}{m_b} \) and \( g_A = 1 - \frac{m_\pi}{m_b} \).

For the evaluation of the form factors \( f_2^T(0) \) and \( g_2^T(0) \) we employ the relativistic quark-diquark model. In the quasipotential approach the matrix element of a local current \( J_\mu \) is given by [9]

\[
\langle \Xi^{-}(P)|J_\mu|\Xi^{-}_b(Q)\rangle = \int \frac{d^3p d^3q}{(2\pi)^6} \bar{\Psi}_{\Xi^{-}_p}(p) \Gamma_\mu(p, q) \Psi_{\Xi^{-}_q}(q),
\]

where \( P \) and \( Q \) are momenta of the final and initial baryons, respectively, and \( \Gamma_\mu(p, q) \) is the two-particle vertex function. In our case \( \Gamma_\mu(p, q) = \Gamma^{(1)}_{\mu}(p, q) + \Gamma^{(2)}_{\mu}(p, q) \), where:

\[
\Gamma^{(1)}_{\mu}(p, q) = \bar{\psi}_d(p_d) \bar{u}_s(p_s) \gamma_\mu (1 - \gamma_5) u_b(q_b) \psi_d(q_d) (2\pi)^3 \delta(p_d - q_d)
\]

is the vertex function corresponding to the impulse approximation diagram (see Fig. 1 from [8], for the analogous process \( \Lambda_b \to \Lambda\gamma \)), while

\[
\Gamma^{(2)}_{\mu}(p, q) = \bar{\psi}'_d(p_d) \bar{u}_s(p_s) \left[ \gamma_\mu (1 - \gamma_5) \frac{\Lambda_{b}^{-}(k)}{\epsilon_b(k) + \epsilon_b(p_s)} \gamma^0 V(p_d - q_d) + \right.
\]

\[
+ \left. \gamma^0 \gamma_\mu (1 - \gamma_5) \frac{\Lambda_{s}^{-}(k')}{\epsilon_s(k') + \epsilon_s(q_b)} \psi_d(q_d) \right]
\]

is the vertex function corresponding the diagrams (see Fig. 2 from [9], for the analogous process \( \Lambda_b \to \Lambda\gamma \)) with the intermediate negative-energy states which are the consequence of the projection onto the positive-energy states in the quasipotential approach. Here \( \bar{\psi}_d(p) \) is the diquark wave function, \( V(p) \) is the quark-diquark interaction quasipotential, \( k = p_s - \Delta, k' = q_b + \Delta, \Delta = P - Q, \epsilon(p) = \sqrt{m^2 + p^2} \), and

\[
\Lambda^{(-)}(p) = \frac{\epsilon(p) - (m \gamma^0 + \gamma^0 (\gamma p))}{2 \epsilon(p)}
\]

and \( u_q(p) \) are the Dirac bispinors, \( m_q \) and \( m_d \) are the quark and diquark masses, respectively.

The baryon wave functions \( \bar{\Psi}_{\Xi^{-}_p}(q) \) and \( \bar{\Psi}_{\Xi^{-}_p}(p) \) are projected onto the positive-energy states of quarks and boosted to the moving reference frame. Indeed, in the rest frame of the initial baryon \( \Xi^{-} \) the final baryon is moving with the recoil momentum \( P \). Thus we have to boost the wave function of the final baryon \( \Xi^{-} \) to the moving reference frame

\[
\bar{\Psi}_{\Xi^{-}p}(p) = D^{1/2}_q(R_{L_p}^W) D_d(R_{L_p}^W) \Psi_{\Xi^{-}0}(p),
\]

where \( \Psi_{\Xi^{-}0}(p) \equiv \Psi_{\Xi^{-}p}(p) \) is the baryon wave function in the rest frame, \( R^W \) is the Wigner rotation, \( L_P \) is the Lorentz boost from the baryon rest frame to a moving one with the
momentum $\mathbf{P}$, and $D_1^{1/2}(R^W)$ is the rotation matrix of the quark spin, while the rotation matrix for the scalar diquark $D_d(R^W) = 1$.

The baryon $B = \Xi_b$ or $\Xi$ wave functions in the rest frame satisfy the relativistic quasipotential equation of the Schrödinger type

\[
\left( \frac{b^2(M)}{2\mu_R(M)} - \frac{\mathbf{p}^2}{2\mu_R(M)} \right) \Psi_B(\mathbf{p}) = \int \frac{d^3q}{(2\pi)^3} V(\mathbf{p}, \mathbf{q}, M) \Psi_B(\mathbf{q}),
\]

where

\[
\mu_R(M) \equiv \frac{M_B^4 - (m_q^2 - m_d^2)^2}{4M_B^2}
\]

is the relativistic reduced mass, and

\[
b^2(M) = \frac{(M_B^2 - (m_q + m_d)^2)(M_B^2 - (m_q - m_d)^2)}{4M_B^2}
\]

is the relativistic center-of-mass system relative momentum squared on the mass shell. The quark-diquark interaction potential $V(\mathbf{p}, \mathbf{q}, M)$ is constructed from the off-mass-shell scattering amplitude projected on the positive energy states. It includes all spin-dependent and spin-independent relativistic contributions. Its explicit form can be found Ref. [10].

Explicit expressions for the form factors are given in Ref. [9]. Substituting the baryon wave functions which were obtained in the baryon mass calculations [11] we obtain the following values of form factors at $q^2 = 0$

\[
f_2^T(0) = g_2^T(0) = -0.144.\]

We estimate the uncertainties of the calculated values of the form factors to be less than 5%.

### IV. RARE RADIATIVE DECAY RATE

The exclusive rare radiative decay rate $\Xi_b^- \to \Xi^- \gamma$ for the emission of a real photon $q^2 = 0$ is given by

\[
\Gamma = \frac{G_F^2 \alpha_{em}}{64\pi^4} |V_{tb}V_{ts}|^2 m_b^2 C_7^{(0)eff} (m_b)^2 (g_V^2 |f_2^T(0)|^2 + g_A^2 |g_2^T(0)|^2) \left( \frac{m_{\Xi_b}^2 - m_{\Xi^-}^2}{m_{\Xi_b}^-} \right)^3,
\]

where $\alpha_{em} \equiv \frac{e^2}{4\pi}$ is the electromagnetic coupling constant.

Substituting the values of the physical constants summarized in Table I [12] and the calculated values of the form factors given in Table II we get the prediction for the branching fraction

\[
Br(\Xi_b^- \to \Xi^- \gamma) = (0.95 \pm 0.15) \times 10^{-5}.
\]

We compare our result with the previous theoretical predictions [3–6] and the experimental upper limit in Table III. One can see, that the result of the light cone sum rules Ref. [6] is significantly higher than other theoretical predictions and it exceeds the experimental upper limit. Our result is consistent with the values from Refs. [3–5] within theoretical uncertainties and these values are lower than the experimental limit. Thus, the measurement of the rare radiative $\Xi_b^- \to \Xi^- \gamma$ decay branching fractions can discriminate between different approaches.
TABLE I: Values of the physical constants.

| Quantity         | Numerical value                     |
|------------------|-------------------------------------|
| $G_F$            | $1.166 \times 10^{-5}$ GeV$^{-2}$   |
| $\alpha_{\text{em}}(M_W)$ | $1/128$                           |
| $|V_{ts}|$        | $(38.8 \pm 1.1) \times 10^{-3}$    |
| $|V_{tb}|$        | $1.013 \pm 0.030$                  |
| $m_b$ (pole)     | $(4.78 \pm 0.06)$ GeV              |
| $m_s$            | $93^{+11}_{-5}$ MeV                |
| $m_{\Xi^-}$      | $(5797.0 \pm 0.6)$ MeV             |
| $m_{\Xi^-}$      | $(1321.71 \pm 0.07)$ MeV           |
| $M_W$            | $(80.379 \pm 0.012)$ GeV           |
| $M_Z$            | $(91.1876 \pm 0.0021)$ GeV         |
| $m_t$            | $(172.76 \pm 0.30)$ GeV            |
| $\tau_{\Xi_b}$  | $(1.572 \pm 0.040) \times 10^{-12}$s |
| $\hbar$          | $6.582 \times 10^{-22}$ MeV s      |

TABLE II: The calculated values of the effective Wilson coefficient and form factors at $q^2 = 0$.

| Quantity         | Numerical value                     |
|------------------|-------------------------------------|
| $C^{(0)\text{eff}}_T(m_b)$ | $-0.310$                           |
| $f_2^T(0) = g_2^T(0)$ | $-0.144$                           |

V. CONCLUSION

The rare radiative decay $\Xi_b^- \rightarrow \Xi^- \gamma$ is investigated in the framework of the relativistic quark model. First, we give the expression for the effective Hamiltonian and evaluate the relevant Wilson coefficient by solving the system of the renormalization group equations.

TABLE III: Comparison of the theoretical predictions with experimental upper limit for the branching fraction of the $\Xi_b^- \rightarrow \Xi^- \gamma$ decay.

| Reference                      | Predicted value                     |
|--------------------------------|-------------------------------------|
| Light cone sum rules [6]       | $(3.03 \pm 0.10) \times 10^{-4}$    |
| SU(3) flavor symmetry [3]      | $(1.23 \pm 0.64) \times 10^{-5}$    |
| Light-cone QCD sum rules in full theory [5] | $1.08^{+0.63}_{-0.49} \times 10^{-5}$ |
| Light-front quark model [4]    | $(1.1 \pm 0.1) \times 10^{-5}$      |
| This paper                     | $(0.95 \pm 0.15) \times 10^{-5}$    |
| Experiment [1]                 | $< 1.3 \times 10^{-5}$              |
Then the quasipotential approach and relativistic quark-diquark picture are employed for the calculation of the form factors parameterizing the hadronic matrix elements of the weak current. The baryon wave functions are obtained with the account of all relativistic effects including transformation from rest to the moving reference frame and contributions of the intermediate negative-energy states. The form factors of this decay are expressed as the overlap integrals of the initial and final baryon wave functions [9]. Using the calculated values of form factors at $q^2 = 0$ (since the real photon is emitted) we obtained the value of the branching ratio. Our result is consistent with the predictions of the theoretical approaches in Refs. [3–5] and is below the experimental upper limit set by LHCb [1], while the branching ratio obtained in Ref. [6] is significantly higher and contradicts the experimental limit. Therefore the exact experimental value is needed to make a final comparison.

Acknowledgments

We are grateful to D. Ebert for valuable discussions.

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