第 33 卷 第 1 期 原 子 核 物 理 评 论 Vol. 33, No. 1
2016 年 3 月 Nuclear Physics Review Mar., 2016

Article ID: 1007-4627(2016) 01-0030-06

Radiative Decays of the Bottom Strange Mesons in a Relativistic Quark Model

WANG Wenjun[1], JIA Duojie[1], CHEN Dianyong[2]

(*1. Institute of Theoretical Physics, College of Physics and Electronic Engineering, Northwest Normal University, Lanzhou 730070, China;*
2. Institute of Modern Physics, Chinese Academy of Sciences, Lanzhou 730000, China)

Abstract: We systematically study the radiative transitions of bottom-strange mesons in the framework of the relativistic constituent quark model. The partial widths of the E1 and M1 decays are predicted. The results predict that most of E1 decay widths are several keV and most of M1 decay widths are less than 1 keV,which give a roadmap of searching for the higher bottom-strange mesons via radiative decays. The experimental searches by the forthcoming Belle II can in turn provide further tests to our result in the present work.

Key words: bottom-strange meson; radiative decay; relativistic constituent quark model

CLC number: O572.33$^+$9 **Document code:** A **DOI:** 10.11804/NuclPhysRev.33.01.030

1 Introduction

In recent decades, much progress on experimental observations of new hadron states has been achieved. A growing number of orbital and radial excited hadrons have been discovered[1], which makes the hadron family lengthy. On the other hand, understanding of the properties of the newly observed hadrons can enlarge our knowledge about quantum chromodynamics(QCD), especially about QCD at low energy. While the non-pertubative nature of QCD at low energy make it difficult to study the hadron properties from the first principle, we rely usually on some QCD-inspired quark models to depict hadron spectra and to investigate hadron decays.

Of all mesons, the heavy-light mesons family are particular interesting since the spin of the heavy quark in a heavy-light meson conserves in the heavy quark limit. This extra symmetry indicates that the wave function of a heavy-light meson has to be independent of the flavor and spin of the heavy quarks[2–5]. In view of the relativistic effects, the mass spectra of heavy-light mesons are well described by the relativistic constituent quark model proposed in Refs. [6–8].

Besides the mass spectra, the decay behaviors of the heavy-light mesons provide further restrict to the wave functions of mesons, which is also directly related to the interactions of both heavy and light quarks. Thus, the systematical investigation on the heavy-light meson decays provides a crucial test to our understanding of the interquark interactions in mesons. As the dominant decay modes of the higher excited bottom strange mesons, the strong decays had been studied via the quark pair creation mechanism in relativistic constituent quark model[9]. Compared to the strong decays, the mechanism of the radiative decays is well understood, thereby the Hamiltonian of photon emission interaction could be constructed. Furthermore, the radiative transitions of mesons are sensitive to the inner quark structure of mesons themselves and therefore, the study of the radiative decays of the heavy-light mesons will be helpful in determining the inner properties of a meson, such as, quantum numbers. In addition, the radiative decays could be a valuable tool of searching for the missing higher excited states.

The mass spectra of the heavy-light meson spectra had been well reproduced[7–8]. With the numerically solved wave functions, one can further study the radiative decays of the heavy-light mesons. Compared to the charmed and charmed-strange mesons, the theoretical estimations for the radiative decays of the bottom-

Received date: 21 Dec. 2015; **Revised date:** 18 Jan. 2016
Foundation item: National Natural Science Foundation of China (11265014, 11375240)
Biography: WANG Wenjun(1989–), male, Gansu province, master, working on particle and nuclear physics; E-mail: wenjun19890128@sina.com
Corresponding author: JIA Duojie, E-mail: jiadj@nwnu.edu.cn.

strange mesons are rare. In the present work, we perform a systematical study of the radiative decays of bottom strange mesons in the framework of relativistic constituent quark model[7–8]. With this investigation, we can provide a roadmap of searching for higher bottom-strange mesons via radiative decays.

This paper is organized as follows, after introduction, a brief review of the relativistic quark model and the formula for the radiative decays are presented in Section 2, the numerical results and some discussions are given in Section 3.

2 Relativistic constitute quark model and radiative decays

Firstly, we present a brief review of the relativistic constituent quark model[7–8]. In view of the relativistic effects, the bounding system of the quark and antiquark is described by the Hamiltonian in the form,

$$\mathcal{H}=\sqrt{p^2+m_1^2}+\sqrt{p^2+m_2^2}+V_{\rm eff}(r)\ , \tag{1}$$

where p is the momentum of the (heavy) quark or the (light) anti-quark in the center-of-mass frame of the meson, m_1 and m_2 are the masses of them, respectively. The effective potential $V_{\rm eff}(r)$ includes the spin-independent confinement potential $V_{\rm Con}(r)$ and spin dependent part composed of the hyperfine interaction, $V_{\rm hyp}(r)$, and the spin-orbital interaction terms, $V_{\rm SO}(r)$. Explicitly, these potentials are given by

$$\begin{aligned} V_{\rm Con}(r)&=-\frac{4}{3}\frac{\alpha_{\rm s}(r)}{r}+br+c\ ,\\ V_{\rm hyp}(r)&=f(r)\boldsymbol{s}_1\cdot\boldsymbol{s}_2+g(r)\left(3(\boldsymbol{s}_1\cdot\hat{r})(\boldsymbol{s}_2\cdot\hat{r})-\boldsymbol{s}_1\cdot\boldsymbol{s}_2\right)\ ,\\ V_{\rm SO}(r)&=h_1(r)\boldsymbol{s}_1\cdot\boldsymbol{L}+h_2(r)\boldsymbol{s}_2\cdot\boldsymbol{L}\ , \end{aligned}\tag{2}$$

where the form of confinement potential is consistent with the calculation of Lattice QCD[10] and the parameter c is an additional constant fitted by the ground state of the spectrum. The r-dependent functions, $f(r)$, $g(r)$, $h_1(r)$ and $h_2(r)$, in the hyperfine and spin-orbital interactions are defined by $\alpha_{\rm s}(r)$ and $\mathrm{d}V_{\rm Con}(r)/\mathrm{d}r$, which read[7],

$$f(r)=\frac{32\pi}{9m_1\tilde{m}_{2a}}\alpha_{\rm s}(r)\delta_\sigma(r)\ , \tag{3}$$

$$g(r)=\frac{4}{3}\frac{\alpha_{\rm s}(r)}{m_1\tilde{m}_{2b}}\frac{1}{r^3}\ , \tag{4}$$

$$h_1(r)=\left[\frac{4}{3}\frac{\alpha_{\rm s}(r)}{r^3}\left(\frac{1}{m_1}+\frac{1}{\tilde{m}_{2c}}\right)-\frac{1}{2r}\frac{\partial V_{\rm Con}(r)}{\partial r}\frac{1}{m_1}\right]\frac{1}{m_1}\ , \tag{5}$$

$$h_2(r)=\left[\frac{4}{3}\frac{\alpha_{\rm s}(r)}{r^3}\left(\frac{1}{m_1}+\frac{1}{\tilde{m}_{2c}}\right)\frac{1}{\tilde{m}_{2c}}-\frac{1}{2r}\frac{\partial V(r)}{\partial r}\frac{1}{(\tilde{m}_{2d})^2}\right]\ , \tag{6}$$

where $\tilde{m}_{2i}$ $(i=a,b,c,d)$ stand for the (dressed) masses of the light antiquark, depending on the quark-antiquark systems considered.

Instead of using the harmonic oscillator basis[6], one can adopt the spherical Bessel functions basis $J_l(kr)$ being restricted in a finite space $r<L$ as a complete basis to expand the bound-state wave function $\psi(r)$,

$$\begin{aligned}\psi(\boldsymbol{r})&=\Phi_{nl}Y_{lm}(\hat{r})=\frac{u_{nl}}{r}Y_{lm}(\hat{r}),\\ u_{nl}&=\sum_{i=1}^{N}c_i^{(nl)}\frac{a_ir}{L}J_l\left(\frac{a_ir}{L}\right)\ ,\end{aligned}\tag{7}$$

where c_i^{nl} are the expansion coefficients, a_i is the i-th root of the spherical Bessel function. For the first step, one can discard the hyperfine and spin-orbital interactions in the Hamiltonian, and solve the Schrödinger equation,

$$H^0|\psi\rangle=E^0|\psi\rangle\ ,$$

with

$$H^0=\sqrt{p^2+m_1^2}+\sqrt{p^2+m_2^2}+V_{\rm Con}(r)\ . \tag{8}$$

It is not hard to transform the above eigenvalue equation into a system of linear equations for the coefficents c_i^{nl}[7–8].

For the second step, one can include the spin-orbital and hyperfine interaction as a perturbative terms. It should be noticed that the tensor term in the hyperfine interaction $V_{\rm hyp}(r)$ does not conserve the orbital angular momentum, thereby giving rise to the mixing between the states with different orbital angular momenta but with the same spin and total angular momenta, *i.e.*, the mixing between ^3L_J and $^3L'_J$. As for the spin-orbit interaction, $V_{\rm SO}(r)$, it does not conserve the total spins of the quark and anti-quark, therefore the states with different total spin quantum numbers but with the same orbital and total angular momenta, *i.e.*, ^1L_J and ^3L_J, will be mixed due to this spin-orbit interaction. Such a kind of mixing can only exist in the mesons compose of different kinds of quark and anti-quark, in which the C-parity of the mesons cannot be defined.

The radiative transition Hamiltonian for mesons can be constructed from the photon-quark interactions, which is,

$$H_{\rm EM}=\frac{e_{\rm q}}{m_{\rm q}}\left(-{\rm i}p'\cdot\epsilon_\gamma^*+\frac{1}{2}\sigma\cdot q\times\epsilon_\gamma^*\right){\rm e}^{-{\rm i}q\cdot r_{\rm q}}+(q\rightarrow\bar{Q})\ , \tag{9}$$

where p' and q are the momentum of quark and photon, respectively, and $e_{\rm q}$ and $m_{\rm q}$ the charge and the mass of the quark. The first term in the bracket is corresponding to E1 transition, which is independent of the quark spin, so this term conserves the spin of the meson, but changes the orbital angular momentum by ±1. After some analytical simplifications, the E1 transitions between $B_{\rm s}$ mesons has the form[11–12],

$$\Gamma^{\rm E1}_{\rm A\rightarrow B}=\frac{4\alpha}{3}\langle e_{\rm M}\rangle^2\omega^3C_{\rm fi}\delta_{\rm SS'}\delta_{\rm LL'+1}|\langle\psi_{\rm B}(r)|r|\psi_{\rm A}(r)\rangle|^2\ , \tag{10}$$

where α is the fine-structure constant, ω is the photon's energy and $e_{\rm M}=(e_{\rm q}m_{\rm Q}-e_{\bar{\rm Q}}m_{\rm q})/(m_{\bar{\rm Q}}+m_{\rm q})$. In the present work we chose $\bar{\rm Q}$ and q as the $\bar{\rm b}$ quark and s quark, respectively. The angular momentum matrix element, $C_{\rm fi}$, is defined as,

$$C_{\rm fi}=\max(L_{\rm A},L_{\rm B})(2J_{\rm B}+1)\left\{\begin{matrix}L_{\rm B} & J_{\rm B} & S\\ J_{\rm A} & L_{\rm A} & 1\end{matrix}\right\}^2\ , \tag{11}$$

where $\{...\}$ is a 6-j symbol.

The second term in the first bracket of Eq. (9) is corresponding to magnetic dipole (M1) transition. The operator σ will change the spin of the quark. The partial width of the magnetic dipole transitions is[12],

$$\Gamma^{\rm M1}({\rm A}\rightarrow{\rm B})=\frac{\alpha\omega^3}{3}\frac{2J_{\rm B}+1}{2J_{\rm A}+1}\delta_{\rm S,S'\pm1}\times \left|\langle\psi_{\rm B}(r)|f_{\rm M1}(r)|\psi_{\rm A}(r)\rangle\right|^2\ , \tag{12}$$

with

$$f_{\rm M1}(r)=\frac{e_{\bar{\rm Q}}}{m_{\bar{\rm Q}}}J_0\left(\frac{\omega m_{\rm q}r}{m_{\bar{\rm Q}}+m_{\rm q}}\right)+\frac{e_{\rm q}}{m_{\rm q}}J_0\left(\frac{\omega m_{\bar{\rm Q}}r}{m_{\bar{\rm Q}}+m_{\rm q}}\right)\ , \tag{13}$$

where $J_0(x)$ is the spherical Bessel function.

3 Numerical results and discussion

All of the involved parameters in the relativistic constituent quark model are listed in Table 1. With these parameters, one can obtain the mass spectra of the bottom-strange mesons. At present, only four bottom-strange mesons have been observed experimentally, which are $B_{\rm s}(1^1S_0)$, $B_{\rm s}(1^3S_1)$, $B_{\rm s}(1^3P_1)$ and $B_{\rm s}(1^3P_2)$, respectively. In Table 2, we present our numerical results for mass spectra of the bottom strange mesons. The experimentally measured values and some other theoretical estimations are also listed for comparison. One can find that the theoretical results in the present relativistic constituent quark model reproduce the experimental data better than the other models, in particular, the discrepancies of the theoretical results and the experimentally measured values are less than 20 and 1 MeV for the S wave and P wave states, respectively.

Table 1 The parameters that this work involved.

Parameter	Value/GeV	Parameter	Value	Parameter	Value	Parameter	Value/GeV
$m_{\rm b}$	4.99	$m_{\rm s}$	0.30 GeV	b	0.16 GeV^2	c	−0.28
σ_0	1.80	S_0	1.55	$\widetilde{m}_{2a}$	0.679 GeV	$\widetilde{m}_{2b}$	0.747
$\widetilde{m}_{2c}$	0.747	$\widetilde{m}_{2d}$	0.889 GeV	L	10 fm		

Table 2 The theoretical evaluation of the bottom-strange meson masses (in unit of MeV). For comparison, the experimentally measured values and some other theoretical results are also listed.

State	J^P	Exp.	Present work	GI model[6]	Ref. [13]	Ref. [14]	Ref. [21]	Ref. [22]	Ref. [23]
$B_{\rm s}(1^1S_0)$	0^-	5 366.77±0.24	5 350	5 390	5 366	−−	5 378	5 372	5 373
$B_{\rm s}(1^3S_1)$	1^-	$5\,415.4^{+2.4}_{-2.1}$	5 400	5 450	5 412	−−	5 440	5 414	5 421
$B_{\rm s}(1^1P_1)$	1^+	5 829.4±0.7	5 820	−−	5 720	5 834	−−	5 865	5 842
$B_{\rm s}(1^3P_2)$	2^+	5 839.7±0.6	5 840	5 880	−−	5 846	−−	5 842	5 820
$B_{\rm s}(1^3D_1)$	1^-	−−	6 090	−−	−−	−−	−−	6 209	6 127
$B_{\rm s}(1^3P_0)$	0^+	−−	5 720	−−	−−	−−	−−	5 833	5 804
$B_{\rm s}(1^3P_1)$	1^+	−−	5 750	−−	−−	−−	−−	5 831	5 805
$B_{\rm s}(2^3P_2)$	2^+	−−	6 260	−−	−−	−−	−−	6 359	6 292
$B_{\rm s}(2^1S_0)$	0^-	−−	5 890	5 980	−−	−−	−−	5 976	5 985
$B_s(2^3S_1)$	1^-	−−	5 940	6 010	−−	−−	−−	5 992	6 019
$B_{\rm s}(2^3D_1)$	1^-	−−	6 410	−−	−−	−−	−−	6 629	−−

Table 2 (Continued)

$B_s(2^3P_0)$	0^+	--	6080	--	--	--	--	6318	6292
$B_s(2^3P_1)$	1^+	--	6150	--	--	--	--	6345	6278
$B_s(2^1P_1)$	1^+	--	6240	--	--	--	--	6321	6296
$B_s(3^1S_0)$	0^-	--	6330	--	--	--	--	6467	6421
$B_s(3^3S_1)$	1^-	--	6380	--	--	--	--	6475	6449
$B_s(1D_2')$	2^-	--	6100	--	--	--	--	--	--
$B_s(1D_2)$	2^-	--	6110	--	--	--	--	--	--
$B_s(2D_2')$	2^-	--	6430	--	--	--	--	--	--
$B_s(2D_2)$	2^-	--	6440	--	--	--	--	--	--

With these predicted masses and the wave functions estimated from the present model, we can perform the calculation of the radiative transitions between the bottom-strange mesons. It should be noticed that no experimental measurement for the radiative transitions between bottom-strange mesons is available yet and also, before the systematical estimations of the radiative transitions in the present work the theoretical predictions are very rare, which makes the present results are valuable for further experimental measurements and for the incoming theoretical investigations on this issue.

The theoretical estimations for the partial widths of the E1 transition are given in Table 3. Most of the partial widths of the E1 transitions are less than 1 keV for the $S\to P\gamma$ transitions, except for the $B_s(3^1S_0)\to B_s(2P_1')\gamma$, $B_s(3^3S_1)\to B_s(2P_1')\gamma$ and $B_s(3^3S_1)\to B_s(2P_1)\gamma$, which are 3.50, 1.80 and 2.43 keV, respectively. As for the $P\to S\gamma$ transitions, most of the partial widths are several keV . As for $D\to P\gamma$ transitions, the partial widths of several channels, such as $B_s(1^3D_1)\to B_s(1^3P_0)\gamma$, $B_s(1^1D_2)\to B_s(1P_1')\gamma$, $B_s(2^1D_2)\to B_s(2P_1')\gamma$, $B_s(2^3D_2)\to B_s(2P_1')\gamma$, are larger than 5 keV. We notice that in Ref. [13] the partial widths of $B_s(1P_1)\to B_s(1^1S_0)\gamma$ and $B_s(1P_1)\to B_s(1^1S_0)\gamma$ are estimated to be $3.2\sim15.8$ keV and $0.3\sim6.1$ keV, which is consistent with our results. In Ref. [16], the partial width for $B_s(1P_1)\to B_s(1^1S_0)\gamma$ is predicted to be $2.01\sim2.67$ keV, in which the $B_s(1P_1)$ was considered as a BK molecule states. In Ref. [15], the predicted partial widths of $B_s(1P_1')\to B_s(1^3S_1)\gamma$ and $B_s(1P_1')\to B_s(1^3S_1)\gamma$ were 56.9 and 39.1 keV in the framework of heavy quark effective theory, which are about 10 times larger than our present results.

Table 3　The partial widths of the E1 transition for B_s mesons, where the partial width below 0.01 keV is marked as ~ 0. The values in bracket are the theoretical predicted masses in unit of MeV.

Initial	Final	Width/keV	Initial	Final	Width/keV
$B_s(2^1S_0)$	$B_s(1P_1')$	0.33	$B_s(2P_1')$	$B_s(1^3S_1)$	65.57
	$B_s(1P_1)$	0.10		$B_s(1^1S_0)$	~ 0
$B_s(2^3S_1)$	$B_s(1^3P_2)$	0.15	$B_s(2P_1)$	$B_s(1^1S_0)$	~ 0
	$B_s(1^3P_0)$	0.76		$B_s(1^3S_1)$	~ 0
	$B_s(1P_1')$	0.64	$B_s(1^3D_1)$	$B_s(1P_1)$	1.66
	$B_s(1P_1)$	0.16		$B_s(1^3P_0)$	5.69
$B_s(3^1S_0)$	$B_s(1P_1')$	0.54		$B_s(2^3P_0)$	~ 0
	$B_s(1P_1)$	0.36		$B_s(1^3P_2)$	0.09
	$B_s(2P_1')$	3.50		$B_s(1P_1')$	3.31
	$B_s(2P_1)$	0.44	$B_s(2^3D_1)$	$B_s(2^3P_0)$	6.13
$B_s(3^3S_1)$	$B_s(2^3P_2)$	0.58		$B_s(1P_1')$	~ 0
	$B_s(2^3P_0)$	1.80		$B_s(1P_1)$	~ 0
	$B_s(2P_1')$	2.43		$B_s(1^3P_0)$	~ 0
	$B_s(2P_1)$	0.55		$B_s(2P_1')$	2.25
	$B_s(1^3P_2)$	0.24		$B_s(2P_1)$	0.63
	$B_s(1^3P_0)$	0.09		$B_s(1^3P_2)$	~ 0
	$B_s(1P_1')$	0.23		$B_s(2^3P_2)$	~ 0
	$B_s(1P_1)$	0.16	$B_s(1^1D_2)$	$B_s(1P_1')$	7.08
$B_s(1^3P_2)$	$B_s(1^3S_1)$	6.36		$B_s(1P_1)$	3.7

Table 3 (Continued)

$B_s(1P_1')$	$B_s(1^3S_1)$	3.20	$B_s(2^1D_2)$	$B_s(2P_1)$	1.84
	$B_s(1^1S_0)$	4.78		$B_s(2P_1')$	5.62
$B_s(1P_1)$	$B_s(1^1S_0)$	7.75		$B_s(1P_1')$	~ 0
	$B_s(1^3S_1)$	5.53		$B_{s1}(1P_1)$	~ 0
$B_s(1^3P_0)$	$B_s(1^3S_1)$	2.45	$B_s(1^3D_2)$	$B_s(1^3P_2)$	0.89
$B_s(2^3P_2)$	$B_s(2^3S_1)$	5.09		$B_s(1P_1')$	6.5
	$B_s(1^3S_1)$	~ 0		$B_s(1P_1)$	3.33
$B_s(2P_1')$	$B_s(2^3S_1)$	1.44	$B_s(2^3D_2)$	$B_s(1P_1')$	~ 0
	$B_s(2^1S_0)$	2.73		$B_s(1P_1)$	~ 0
$B_s(2P_1)$	$B_s(2^1S_0)$	6.66		$B_s(1^3P_2)$	~ 0
	$B_s(2^3S_1)$	4.20		$B_s(2^3P_2)$	0.4
$B_s(2^3P_0)$	$B_s(2^3S_1)$	0.43		$B_s(2P_1')$	5.06
	$B_s(1^3S_1)$	~ 0		$B_s(2P_1)$	1.58

As for the M1 transitions, we list our results in Table 4. Most of the predicted partial widths are less than 1 keV, except for those of $B_s(1P_1) \to B_s(1P_1')\gamma$, $B_s(1^3P_2) \to B_s(1P_1')\gamma$, $B_s(2P_1) \to B_s(2P_1')\gamma$, $B_s(2^3P_2) \to B_s(2P_1')\gamma$, which are estimated to be 1.16, 2.46, 2.45 and 4.48 keV, respectively. In addition, the transition between $B_s(1^3S_1)$ and $B_s(1^1S_1)$ should be a interesting channel, since both states have been well established experimentally. Theoretically, the partial decay width of $B_s(1^3S_1) \to B_s(1^1S_1)\gamma$ had been predicted to be located in 0.05 to 0.22 keV[17–20], which are smaller than our results.

Table 4 The partial widths of the E1 transition for B_s mesons, where the partial width below 0.01 keV is marked as ~ 0. The values in the bracket are the theoretical predicted masses in unit of MeV.

Initial	Final	Width/keV	Initial	Final	Width/keV
$B_s(1^3S_1)$	$B_s(1^1S_0)$	0.42	$B_s(2^3P_2)$	$B_s(2P_1')$	4.48
$B_s(2^1S_0)$	$B_s(1^3S_1)$	0.07		$B_s(2P_1)$	0.03
$B_s(2^3S_1)$	$B_s(2^1S_0)$	0.42		$B_s(1P_1')$	0.08
	$B_s(1^1S_0)$	0.08		$B_s(1P_1)$	0.03
$B_s(3^3S_1)$	$B_s(3^1S_0)$	0.42	$B_s(2^3D_1)$	$B_s(1D_2)$	~ 0
	$B_s(2^1S_0)$	0.05		$B_s(1D_2')$	~ 0
	$B_s(1^1S_0)$	0.02	$B_s(2D_2)$	$B_s(1^3D_1)$	~ 0
$B_s(3^1S_0)$	$B_s(2^3S_1)$	0.07		$B_s(1D_2')$	0.01
	$B_s(1^3S_1)$	0.13		$B_s(1D_2)$	~ 0
$B_s(1P_1)$	$B_s(1P_1')$	1.16		$B_s(2D_2')$	~ 0
	$B_s(1^3P_0)$	0.88		$B_s(2^3D_1)$	0.04
$B_s(1^3P_2)$	$B_s(1P_1')$	2.46	$B_s(2D_2')$	$B_s(1D_2)$	~ 0
	$B_s(1P_1)$	0.01		$B_s(1D_2')$	~ 0
$B_s(2P_1)$	$B_s(2P_1')$	2.45		$B_s(1^3D_1)$	~ 0
	$B_s(1P_1')$	0.05		$B_s(2^3D_1)$	~ 0
	$B_s(1P_1)$	0.02	$B_s(1D_2)$	$B_s(1^3D_1)$	0.01
	$B_s(1^3P_2)$	0.02		$B_s(1D_2')$	~ 0
	$B_s(1^3P_0)$	0.03			

To summarize, we perform, in the present work, a systematical investigation on the radiative decays of the bottom-strange mesons in the framework of relativistic constituent quark model, where both E1 and M1 transitions are discussed for the first time. The results obtained in this work should be useful for further experimental measurements and the incoming theoretical investigations on this issue. In the forthcoming Belle II[24], a large quantity of excited bottom-strange mesons could be produced with an updated center-of-mass energy, which could provide a good platform of measuring the radiative transitions of the bottom-strange mesons. These measurement could provide a crucial test of our present results.

References:

[1] OLIVE K A, AGASHE K, AMSLER C, *et al* (Particle Data Group Collaboration). Chin Phys C, 2014, **38**: 090001.
[2] KAYMAKCALAN O, RAJEEV S, SCHECHTER J. Phys Rev D, 1984, **30**: 594.
[3] OH Y S , SONG T, LEE S H. Phys Rev C, 2001, **63**: 034901.
[4] CASALBUONI R, DEANDREA A, BARTOLOMEO N D, *et al.* Phys Rept, 1997, **281**: 145.
[5] COLANGELO P, FAZIO F DE , PHAM T N. Phys Lett B, 2002, **542**: 71.
[6] GODEFREY S, ISGUR N. Phys Rev D, 1985, **32**: 189.
[7] LIU Jingbin, YANG Maozhi. JHEP, 2014, **1407**: 106.
[8] LIU Jingbin, YANG Maozhi. Phys Rev D. 2015, **91**: 094004.
[9] SUN Yuan, SONG Qintao, CHEN Dianyong, *et al.* Phys Rev D, 2014, **89**: 054026.
[10] BALI G S. Phys Rept, 2001, **343**: 1.
[11] KWONG W, ROSNER J L. Phys Rev D, 1988, **38**: 279.
[12] NOVIKOV V A, OKUN L B, SHIFMAN M A, *et al.* Phys Rept, 1978, **41**: 1.
[13] WANG Zhigang. Commun Theor Phys, 2009, **52**: 91.
[14] EICHTEN E J, HILL C T, QUIGG C. Phys Rev Lett, 1993, **71**: 4116.
[15] BARDEEN W A, EICHTEN E J, HILL C T. Phys Rev D, 2003, **68**: 054024.
[16] FAESSLER A, GUTSCHE T, LYUBOVITSKIJ V E. *et al.* Phys Rev D, 2008, **77**: 114013.
[17] EBERT D, FAUSTOV R N, GALKIN V O. Phys Lett B, 2002, **537**: 241.
[18] ORSLAND A H, HOGAASEN H. Eur Phys J C, 1999, **9**: 503.
[19] ZHU Shi Lin, HWANG W Y P, YANG Z S. Mod Phys Lett A, 1997, **12**: 3027.
[20] JENA S N, PANDA P, SAHU K P. J Phys G, 2001, **27**: 1519.
[21] RUJULA A DE , GEORGI H, GLASHOW S L. Phys Rev D, 1975, **12**: 147.
[22] EBERT D, FAUSTOV R N, GALKIN V O. Eur Phys J C, 2010, **66**: 197.
[23] PIERRO M DI, EICHTEN E. Phys Rev D, 2001, **64**: 114004.
[24] WANG Boqun (the Belle II Collaboration), arXiv: 1511. 09434. 094016.

相对论夸克模型中底奇异介子的辐射衰变

王文君[1]，贾多杰[1]，陈殿勇[2]

(1. 西北师范大学物理与电子工程学院，兰州 730070;
2. 中国科学院近代物理研究所，兰州 730000)

摘要: 系统地研究了相对论夸克模型框架下的底奇异介子的辐射衰变。给出了底奇异介子 E1 和 M1 辐射衰变分宽度。这些结果表明，大多数电偶极辐射衰变宽带能达到数个 keV，大多数磁偶极辐射衰变宽度小于 1 个 keV，这为实验上通过辐射衰变寻找底奇异介子提供了理论依据。建设中的 Belle II 上的实验可以进一步验证我们的结果。

关键词: 底奇异介子；辐射衰变；相对论夸克模型

收稿日期: 2015-12-21； **修改日期:** 2016-01-18
基金项目: 国家自然科学基金项目(11265014, 11375240)
通信作者: 贾多杰，E-mail：jiadj@nwnu.edu.cn。