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Cabibbo-suppressed decays of the Ω_c^0—Feedback to the Ξ_c^+ lifetime

A. Babić, B. Guberina, B. Melić, and H. Štefančić

Theoretical Physics Division, Rudjer Bošković Institute, P.O. Box 180,HR-10002, Zagreb, Croatia

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We investigate a possible background of the type $\Omega_c^0 \to \Xi_c^+ \pi^-$ to the CLEO Ξ_c^+ lifetime measurement. This decay mode may lead to an overestimate of the Ξ_c^+ decay length and, therefore, increase the measured Ξ_c^+ lifetime. The branching ratio $\Gamma(\Omega_c^0 \to \Xi_c^+ \pi^-)/\Gamma(\Omega_c^0 \to \Omega^- \pi^+)$ is analyzed in the framework of the pole model and the modified current algebra. We find that the $\Omega_c^0 \to \Xi_c^+ \pi^-$ decay mode could not generate a substantial systematic error in the Ξ_c^+ lifetime measurement.

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The lifetime measurements of charmed baryons are well known [1] to be very important in estimating and disentangling the different preasymptotic effects in the decays of charmed hadrons. The preasymptotic effects [2], like the inclusion of soft degrees of freedom (light quarks, gluons) generate nonperturbative power corrections, e.g., the destructive and/or constructive Pauli interference, and the W-exchange contribution, producing the diversity of lifetimes of charmed mesons [3–5] and baryons [6,7], which would, otherwise, be all equal in the asymptotic limit of infinitely heavy quark mass[1].

Inclusive hadronic decay rates and lifetimes were calculated a long time ago [3–7] by summing over all possible channels and integrating over some range of energies. A "practical" version of the OPE is used in calculations, i.e., it is assumed that the coefficient functions can be found perturbatively and all nonperturbative effects reside in matrix elements. In real world, however, there are nonperturbative effects even at short distances, and the matrix elements are subject to perturbative corrections too.

Suprisingly enough, the theory works rather well, even in the charmed hadron sector, although the expansion parameter $\sqrt{\mu_G^2(D)/m_c^2} \simeq 0.5$ is not really small (the corresponding parameter in beauty decays is $\sqrt{\mu_G^2(B)/m_b^2} \simeq 0.13$).

A systematic study of charmed baryon decays was performed a few years ago [8], with good agreement between theory and experiment. The theoretical predictions were rather stable to the uncertainties in the wave functions of heavy baryons and/or to the choice of the renormalization/factorization scale, except in the case of the Ξ_c^+ charmed baryon. It was not clear if the peculiar behavior of the Ξ_c^+ was a pure coincidence due to the wild cancellation of different preasymptotic effects, or some deeper understanding was missing. The theoretical result, $\tau(\Xi_c^+)_{\rm th} = 0.27$ ps, for $m_c = 1.35$ GeV, $\Lambda_{\rm QCD} = 300$ MeV, had to be compared with the experimental value, $\tau(\Xi_c^+)_{\rm exp} = (0.35 \pm 0.07)$ ps. The difference, at that time, was not so significant that one would have had to worry. However, it was clear that future more precise measurements could disturb an idyllic concordance between theory and experiment.

Figure 1 shows the results of Ξ_c^+ lifetime experiments performed up to now. One can see that two new measurements with significantly improved accuracy, FOCUS [9] and CLEO [10], are above the previous world average 1σ margin, in the case of CLEO, even above the 2σ margin. By including these two new measurements the average has changed from 0.33 ps to 0.442 ps. In particular, FOCUS precisely measured the charmed-strange baryon Ξ_c^+ lifetime as

$$\tau(\Xi_c^+) = 0.439 \pm 0.022 \pm 0.009 \text{ ps}. \tag{1}$$

In the FOCUS spectrometer, which is well suited to reconstruct short-lived charmed decays, the charmed particles are produced as the product of the interaction between high energy photons with $\langle E \rangle \simeq 180$ GeV in a segmented BeO target and an excellent vertex separation between the production and decay vertices is provided by two silicon vertex detectors.

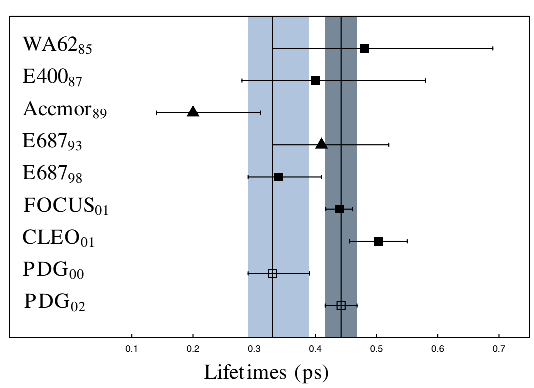

FIG. 1 (color online). Ξ_c^+ lifetime experiments. The left (right) band is the 1σ PDG 2000 [22] (2002 [19]) world average. E687$_{93}$ is excluded from the PDG 2000 value and Accmor from the PDG 2002 value.

[1]It appears astonishing that decay rates of weak and radiative decays are described in terms of few basic quantities, e.g., quark masses, and hadronic expectation values of several leading local operators.

All previous experiments, including that performed by FOCUS, are fixed-target experiments. CLEO performed the only colliding beam experiment. Therefore, it has different systematics and different backgrounds. In spite of the fact that the charmed baryon lifetimes are not measured as precisely as those of charmed mesons, CLEO and SELEX [11] recently measured $\tau(\Lambda_c^+)$ to a precision of 5%. Other charmed baryons (Ξ_c^+, Ξ_c^0, Ω_c^0) are measured with up to an uncertainty of 30%. CLEO's measurement gives

$$\tau(\Xi_c^+) = 0.503 \pm 0.047(\text{stat.}) \pm 0.018(\text{syst.})\ \text{ps}. \qquad (2)$$

This result is obtained using an integrated luminosity of 9.0 fb^{-1} of e^+e^- annihilation data taken with the CLEO IV.V detector at the CESR. The data were taken at energies at and below the $\Upsilon(4S)$ resonance and include $\sim 11 \cdot 10^6$ $e^+e^- \to c\bar{c}$ events. The Ξ_c^+ is reconstructed from the $\Xi^-\pi^+\pi^-$ decay mode. Each Ξ^- is reconstructed from $p\pi^-$. The assumption is that the Ξ_c^+ is produced at the primary event vertex and is not a decay product of another weakly decaying particle, e. g., $\Omega_c^0 \to \Xi_c^+\pi^-$, $\Xi_{cc}^{++} \to \Xi_c^+ + \cdots$.

If Ω_c^0 is produced at the primary event (PE) vertex, travels a certain distance and decays into Ξ_c^+ and π^- (Fig. 2), the production vertex of Ξ_c^+ is misinterpreted to be at the PE vertex and there is an addition (Δ) to the measured proper time:

$$t = \frac{m_{\Xi_c^+}}{c p_{y_{\Xi_c^+}}}(y_{\text{decay}} - y_{\text{production}} + \Delta). \qquad (3)$$

The measured Ξ_c^+ lifetime will be shifted towards a higher value.

The purpose of this letter is to examinate the relevance of the $\Omega_c^0 \to \Xi_c^+\pi^-$ decay mode as a possible source of a systematic error for the Ξ_c^+ lifetime measurement. To this end, we study the ratio of two exclusive decay modes,

$$\eta = \frac{\Gamma(\Omega_c^0 \to \Xi_c^+\pi^-)}{\Gamma(\Omega_c^0 \to \Omega^-\pi^+)} \qquad (4)$$

for the following reasons: the $\Omega_c^0 \to \Omega^-\pi^+$ process is expected to be one of the first and best measured Ω_c^0 exclusive decays in the near future; therefore it is quite convenient to have the contribution of $\Omega_c^0 \to \Xi_c^+\pi^-$ normalized to the rate of $\Omega_c^0 \to \Omega^-\pi^+$ [12]; the $\Omega_c^0 \to \Omega^-\pi^+$ process has a factorizable contribution only, which reduces theoretical uncertainties; uncertainties are further suppressed by considering ratios of exclusive decay widths.

$\Omega_c^0 \to \Xi_c^+\pi^-$ decay mode. In the $\Omega_c^0 \to \Xi_c^+\pi^-$ decay, the decay happens in the light-quark sector and the pion

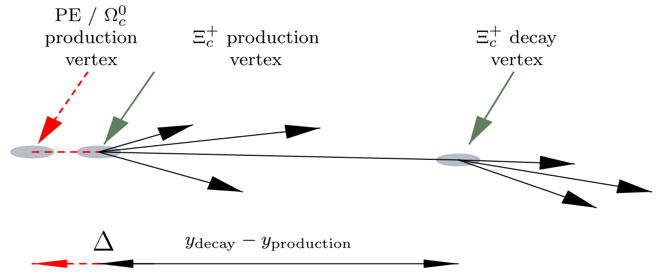

FIG. 2 (color online). Ω_c^0 is produced at the primary event (PE) vertex and decays into Ξ_c^+.

emerges with a momentum of $O(200\ \text{MeV})$ that can be considered 'reasonably" soft. Therefore, there is a similarity between this decay and the hyperon ($\Delta S = 1$) decays for which the soft-pion limit technique with pole corrections was successfully applied [13] with the predictions for the branching fractions within 20% from experimental values. We believe that the soft-pion limit is equally applicable to the $\Omega_c^0 \to \Xi_c^+\pi^-$ decay.

The invariant amplitude for the decay of the initial baryon $B_i(1/2^+)$ to the final baryon $B_f(1/2^+)$ and a pion π^a, $a = 1, 2, 3$, is given by

$$\langle B_f\pi^a|\mathcal{H}_{\text{W}}(0)|B_i\rangle = \mathrm{i}\bar{u}_f(A - B\gamma_5)u_i, \qquad (5)$$

with A and B to be determined using the standard nonleptonic weak hamiltonian

$$\mathcal{H}_{\text{W}} = \sqrt{2}G_F V_{\bar{q}_3q_4}V^*_{\bar{q}_1q_2}(c_-O_- + c_+O_+), \qquad (6)$$

where $O_\pm$ are local 4-quark operators

$$O_\pm = (\bar{q}_{1L}\gamma_\mu q_{2L})(\bar{q}_{3L}\gamma^\mu q_{4L}) \pm (\bar{q}_{3L}\gamma_\mu q_{2L})(\bar{q}_{1L}\gamma^\mu q_{4L}), \qquad (7)$$

with $(\bar{q}_{iL}\gamma_\mu q_{jL}) = \frac{1}{2}\bar{q}_i\gamma_\mu(1-\gamma_5)q_j$, and V's are the elements of the CKM matrix. The Wilson coefficients in the leading logarithmic approximation are given by

$$c_\pm(\mu^2) \cong \left(\frac{\alpha_s(\mu^2)}{\alpha_s(M_{\text{W}}^2)}\right)^{d_\pm/2b}, \qquad (8)$$

where $b = \frac{1}{3}(11N_{\text{c}} - 2n_{\text{f}})$, N_{c} and n_{f} being the number of colors and flavors, respectively. The quantities $d_- = -2d_+ = 8$ are proportional to the anomalous dimensions of the operators O_- and O_+.

In the approach of references [13,14] modified current algebra techniques were applied, i.e., the soft-pion amplitude (commutator term) was corrected for the soft-pion limit. The contribution coming from baryon poles is given as

$$A^{\text{CA}} = A^{\text{soft}} + A^{\text{corr}} = \frac{\sqrt{2}}{f_\pi}\langle B_f|[Q^a, \mathcal{H}_{\text{W}}^{\text{PC}}]|B_i\rangle - \frac{\sqrt{2}}{f_\pi}\sum_{B_n^*(1/2^-)}(m_{B_f} - m_{B_i})\left(\frac{g_A^{B_fB_n^*}b_{B_n^*B_i}}{m_{B_i} - m^*_{B_n}} + \frac{b_{B_fB_n^*}g_A^{B_n^*B_i}}{m_{B_f} - m^*_{B_n}}\right), \qquad (9)$$

$$B^{\text{pole}} = \sum_{B_n(1/2^+)}\frac{\sqrt{2}}{f\pi}\left(\frac{m_{B_f} + m_{B_n}}{m_{B_i} - m_{B_n}}g_A^{B_fB_n}a_{B_nB_i} + \frac{m_{B_i} + m_{B_n}}{m_{B_f} - m_{B_n}}a_{B_fB_n}g_A^{B_nB_i}\right). \qquad (10)$$

In the above equations the S- and P-wave amplitudes are calculated in the framework of the pole model. Using the Lehman-Symanzik-Zimmerman reduction formalism, the pion momentum is taken off shell. The pion field is related to the axial vector curent via PCAC, and a complete set of states is inserted.

In (10), the baryon-baryon weak matrix elements $b^*_{B_jB_i}$ and $a_{B_jB_i}$ are defined as

$$\langle B_j^*(1/2^-)|\mathcal{H}_W^{PV}|B_i\rangle = \mathrm{i}b^*_{B_jB_i}\bar{u}_j u_i, \qquad (11)$$

$$\langle B_j(1/2^+)|\mathcal{H}_W^{PC}|B_i\rangle = a_{B_jB_i}\bar{u}_j\gamma_5 u_i, \qquad (12)$$

and $g_A^{B_iB_j}$ is the axial form-factor, related to the strong $g^{B_iB_jM}$ baryon-baryon-meson coupling through the generalized Goldberger-Treiman relation. The pion decay constant f_π is taken as 0.132 GeV. The weak matrix elements (11) and (12) and the axial form-factors are calculated in the MIT bag model [15].

Concerning the pole resonances, the only flavor structure that can be formed in an intermediate state of the $\Omega_c^0 \to \Xi_c^+\pi^-$ decay is (dsc), (Fig. 3). The main contribution to the S-wave amplitude comes from the commutator term in (17), providing a simple means of summing contributions from all intermediate states in the soft-pion limit. The correction to this term is dominated by $(1/2^-)$ resonances, the lowest one being for our decay $\Xi_c^0(2790)$ (denoted by Ξ_c^{*0}). P-wave amplitudes are dominated by the lowest lying $(1/2^+)$ baryon intermediate states. Since the charmed antitriplet-antitriplet axial form-factors vanish, $g_A^{B_i(\bar{3})B_j(\bar{3})} = 0$, the lowest lying Ξ_c^0 resonance belonging to the charmed baryon antitriplet does not contribute. The main contribution to the P-wave amplitude comes from the $\Xi_c^{\prime 0}$ baryon $(1/2^+)$ state, belonging to the charmed baryon sextet. Therefore we have

$$A^{\mathrm{CA}} = \frac{1}{f_\pi}\langle\Xi_c^0|\mathcal{H}_W^{PC}|\Omega_c^0\rangle + \frac{1}{f_\pi}\frac{m_{\Omega_c^0} - m_{\Xi_c^+}}{m_{\Omega_c^0} - m_{\Xi_c^{0*}}} g_A^{\Xi_c^+\Xi_c^{*0}} b_{\Xi_c^{*0}\Omega_c^0}, \qquad (13)$$

$$B^{\mathrm{pole}} = \frac{1}{f_\pi}\frac{m_{\Xi_c^{\prime 0}} + m_{\Xi_c^+}}{m_{\Omega_c^0} - m_{\Xi_c^{\prime 0}}} g_A^{\Xi_c^+\Xi_c^{\prime 0}} a_{\Xi_c^{\prime 0}\Omega_c^0}. \qquad (14)$$

There is also a factorizable P-wave part of the $\Omega_c^0 \to \Xi_c^+\pi^-$ amplitude, which can be expressed as

$$B^{\mathrm{fact}} = -\frac{G_F}{\sqrt{2}} V_{us}V_{ud} a_1 f_\pi (m_{\Omega_c^0} + m_{\Xi_c^+}) g_A^{\Xi_c^+\Omega_c^0}, \qquad (15)$$

where $a_1 = \frac{1}{3}(2c_+ + c_-)$. The decay rate for $\Omega_c^0 \to \Xi_c^+\pi^-$ is then given by

$$\Gamma(\Omega_c^0 \to \Xi_c^+ + \pi^-) = \frac{|\vec{p}_{\Xi_c^+}|}{4\pi m_{\Omega_c^0}}[|A|^2(E_{\Xi_c^+} + m_{\Xi_c^+}) + |B|^2(E_{\Xi_c^+} - m_{\Xi_c^+})]. \qquad (16)$$

$\Omega_c^0 \to \Omega^-\pi^+$ decay mode. This mode is of the type $B_f(1/2^+) \to B_i(3/2^+) + \pi$ and its invariant amplitude is

$$\mathcal{M} = \mathrm{i}q_\mu \bar{u}_f^\mu (B' - C\gamma_5)u_i. \qquad (17)$$

The expression for the decay rate is

$$\Gamma(\Omega_c^0 \to \Omega^- + \pi^+) = \frac{|\vec{p}_{\Omega^-}|^3 m_{\Omega_c^0}}{6\pi m_{\Omega^-}^2}[|B'|^2(E_{\Omega^-} + m_{\Omega^-}) + |C|^2(E_{\Omega^-} - m_{\Omega^-})]. \qquad (18)$$

The $\Omega_c^0 \to \Omega^-\pi^+$ decay does not receive any pole contributions. There is only a factorizable P-wave amplitude contributing. This decay mode has already been calculated in the literature [16–18] by applying different quark models. We have recalculated it in the MIT bag model in order to have a consistent calculation of the ratio η (4).

Numerical results and discussions. As we have already stated before, all our form-factors, decay constants and matrix elements have been calculated in the MIT bag model. The calculations have been performed using the following parameter set: $\mu = 1$ GeV, $\Lambda_{\mathrm{QCD}} = 200$ MeV, particle masses are taken to be PDG average values [19], and MIT bag model parameters have the same values as in [20].

For the $\Omega_c^0 \to \Xi_c^+\pi^-$ decay mode, we have the S-wave amplitude which is given by the current algebra term and a pole correction of 10%. The P-wave amplitude has a factorizable contribution and a large pole contribution. Note from (16) that the P-wave amplitude is suppressed by a small kinematical factor, making contributions from S- and P-wave amplitudes of the same order of magnitude. The results of the calculation are summarized in Table I. In the $\Omega_c^0 \to \Omega^-\pi^+$ decay the only nonvanishing contribution is from the factorizable part of the P-wave (B') amplitude, the D-wave (C) amplitude is zero (Table II). Finally, the ratio of partial decay rates two

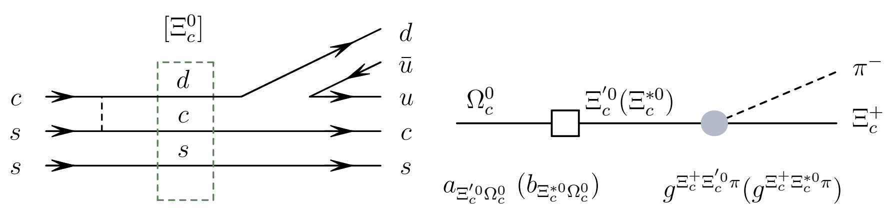

FIG. 3 (color online). Pole diagrams for the $\Omega_c^0 \to \Xi_c^+\pi^-$ decay mode: at the quark level and in terms of effective couplings.

TABLE I. Amplitudes ($\times 10^7$) and width (s^{-1}) for the $\Omega_c^0 \to \Xi_c^+ \pi^-$ decay mode. The invariant amplitude for the decay mode with the spin-1/2 particle in the final state is $\mathcal{M} = i\bar{u}_f(A - B\gamma_5)u_i$, with dimensionless S- and P-wave amplitudes.

$A^{\rm fact}$	$A^{\rm soft}$	$A^{\rm corr}$	$A^{\rm tot}$	$B^{\rm fact}$	$B^{\rm pole}$	$B^{\rm tot}$	$\Gamma(s^{-1})$)
0	2.87	0.25	3.12	7.47	-45.16	-37.69	$4.50 \cdot 10^9$

TABLE II. Amplitudes ($\times 10^7$) and widths (s^{-1}) for the $\Omega_c^0 \to \Omega^- \pi^+$ decay mode. The invariant amplitude for the decay mode with the spin-3/2 particle in the final state is $\mathcal{M} = iq_\mu \bar{u}_f^\mu (B' - C\gamma_5)u_i$, with P- and D-wave amplitudes having units GeV^{-1}.

$B'(\mathrm{GeV}^{-1})$	$C(\mathrm{GeV}^{-1})$	Γ (s^{-1})
13.75	0	2.89×10^{11}

exclusive decay modes of Ω_c^0 considered in this letter is

$$\eta = \frac{\Gamma(\Omega_c^0 \to \Xi_c^+ \pi^-)}{\Gamma(\Omega_c^0 \to \Omega^- \pi^+)} = \frac{2.96 \cdot 10^{-15}\,\mathrm{GeV}}{1.90 \cdot 10^{-13}\,\mathrm{GeV}} = 0.016. \quad (19)$$

The uncertainties, of order 10%, are connected with the scale μ at which the Wilson coefficients are evaluated, whereas the variation of $\Lambda_{\rm QCD}$ from 200 MeV to 300 MeV leads to 15% larger value of η.

The ratio of partial decay rates (19) shows that the branching ratios for the Cabibbo-suppressed decays of Ω_c^0 are at most at the level of a percent. The apparent dilatation of the Ξ_c^+ baryon path due to the described cascade of weak decays from the initially formed Ω_c^0 baryon is quite small and certainly insufficient to explain the discrepancy of the recent Ξ_c^+ lifetime measurements [9,10] and theoretical calculations [8]. This result is altogether not so surprising, although reassuring given that in exclusive decays there is always a possibility of a large pole contribution.

Finally, it is worth mentioning that the improved knowledge on Cabibbo-suppressed decays of singly charmed baryons may have other important implications on the understanding of the Ξ_c^+ lifetime. As shown in [21], it is possible to obtain a model-independent prediction of this lifetime once a reliable estimate of the decay rate of inclusive Cabibbo-suppressed decays of Λ_c^+ is available. Therefore, a more systematic and detailed approach to the Cabibbo-suppressed decays of singly charmed baryons is called for from both the experimental side, as a way of reducing systematic errors, and the theoretical side, as a way of obtaining model-independent predictions of the Ξ_c^+ lifetime.

With the calculated level of the contribution of Cabibbo-suppressed Ω_c^0 decays, it is clear that this form of the systematic error in the determination of the Ξ_c^+ lifetime cannot provide an explanation of the present disaccord between theory and experiment. To achieve agreement, a new layer of theoretical analysis will have to be uncovered and new experimental data will have to be compiled.

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[1] G. Bellini, I. I. Y. Bigi, and P. J. Dornan, Phys. Rep. **289**, 1 (1997); T. Mannel, Acta Phys. Pol. B **32**, 1857 (2001).
[2] B. Blok and M. A. Shifman, in: *Proceedings of the Third Workshop on the Tau-Charm Factory, Marbella, Spain, 1993*, edited by Jasper Kirkby and Rita Kirkby (Editions Frontieres, Gif-sur-Yvette, France, 1994).
[3] B. Guberina, S. Nussinov, R. D. Peccei, and R. Rückl,Phys. Lett. B **89**, 111 (1979).
[4] N. Bilić, B. Guberina, and J. Trampetić, Nucl. Phys. B **248**, 261 (1984).
[5] M. A. Shifman and M. B. Voloshin, Sov. J. Nucl. Phys. **41**, 120 (1985).
[6] M. A. Shifman and M. B. Voloshin, Sov. Phys. JETP **64**, 698 (1986).
[7] B. Guberina, R. Rückl, and J. Trampetić, Z. Phys. C **33**, 297 (1986).
[8] B. Guberina and B. Melić, Eur. Phys. J. C **2**, 697 (1998).
[9] J. M. Link *et al.*, Phys. Lett. B **523**, 53 (2001).
[10] A. H. Mahmood *et al.*, Phys. Rev. D**65**, 031102 (2002).
[11] S. Bianco, F. L. Fabbri, D. Benson, and I. Bigi, Riv. Nuovo. Cim. **26N7-8**, 1 (2003).
[12] Richard Kass, (private communication).
[13] A. Le Yaouanc, O. Pene, J. C. Raynal, and L. Oliver, Nucl. Phys. B **149**, 321 (1979); B. Borasoy and B. R. Holstein, Phys. Rev. D **60**, 054021 (1999).
[14] B. Stech and Q. P. Xu, Z. Phys. C **49**, 491 (1991).
[15] A. Chodos, R. L. Jaffe, K. Johnson, C. B. Thorn, and V. F. Weisskopf, Phys. Rev. D **9**, 3471 (1974); A. Chodos, R. L. Jaffe, K. Johnson, C. B. Thorn, Phys. Rev. D**10**, 2599 (1974).
[16] Q. P. Xu and A. N. Kamal, Phys. Rev. D **46**, 3836 (1992).
[17] J. G. Körner and M. Krämer, Z. Phys. C **55**, 659 (1992).
[18] H. Y. Cheng, Phys. Rev. D **56**, 2799 (1997).
[19] K. Hagiwara *et al.*, Phys. Rev. D**66**, 010001 (2002).
[20] H. Y. Cheng and B. Tseng, Phys. Rev. D **46**, 1042 (1992).
[21] B. Guberina and H. Štefančić, Phys. Rev. D **65**, 114004 (2002).
[22] D. E. Groom *et al.*, Eur. Phys. J. C **15**, 1 (2000).