Rare $\Omega^- \to \Xi^0(1530)\pi^-$ decay in the Skyrme model

G. Duplančić,[1] H. Pašagić,[2] and J. Trampetić[1,3]

[1]*Theoretical Physics Division, Rudjer Bošković Institute, Zagreb, Croatia*
[2]*Faculty of Transport and Traffic Engineering, University of Zagreb, P.O. Box 195, 10000 Zagreb, Croatia*
[3]*Theoretische Physik, Universität München, Theresienstr. 37, 80333 München, Germany*

Rare nonleptonic $\Omega^- \to \Xi^0(1530)\pi^-$ decay branching ratio is estimated by means of the QCD enhanced effective weak Hamiltonian supplemented by the SU(3) Skyrme model used to estimate the nonperturbative matrix elements. Using mean values for experimental input parameters and the Skyrme charge $e = 4.75$ we obtain the rate which is in good agreement with data.

It is well known that the nonleptonic weak decays of baryons can be reasonably well described in the framework of the Standard Model [1,2]. Recently both s-, p-wave nonleptonic hyperon and p-wave Ω^- decay amplitudes were quite successfully reproduced by the SU(3) extended Skyrme model with the QCD enhanced effective weak Hamiltonian [3,4]. The decay amplitudes were described through the current-algebra commutator, the ground-state baryon-pole terms and factorizable contributions. The nonperturbative quantities, i.e., the baryon six dimensional operator matrix elements, were determined using the SU(3) extended Skyrme model. For the s-wave hyperon decay amplitudes correct relative signs and absolute magnitudes were obtained. For the p-waves all relative signs were correct, with their relative magnitudes roughly following the experimental data. One is thus faced with an obvious question: could an analogous approach work equally well for the rare nonleptonic $\Omega^- \to \Xi^0_{3/2^+}(1530)\pi^-$ [5] weak decay? Such a question should be considered in connection with the measurements of the Ω^- lifetime and branching ratios [6]:

$$\tau(\Omega^-)_{\text{exp}} = (82.1 \pm 1.1) \text{ ps} \tag{1}$$

$$\Gamma(\Omega^- \to \Xi^{*0}\pi^-)/\Gamma_{\text{tot}} = \left(6.4 \, {}^{+5.1}_{-2.0}\right) \times 10^{-4}. \tag{2}$$

In this report our goal is to test whether the effective weak Hamiltonian and the SU(3) Skyrme model are able to predict the rare nonleptonic $\Omega^- \to \Xi^{*0}\pi^-$ decay following the method in [3,4]. To this end we shall employ the Standard Model effective Hamiltonian and *the minimal number of couplings concept* of the Skyrme model to estimate the nonperturbative matrix elements of the four-quark operators [3,4,7] for decuplet-decuplet transitions. Throughout this report we use the *arctan ansatz* for the Skyrme profile function $F(r)$ [8], which allows us to calculate the pertinent overlap integrals analytically with accuracy of the order of $\lesssim 1\%$ with respect to the exact numerical results.

The starting point in such an analysis is the effective weak Hamiltonian in the form of the current $\otimes$ current interaction, enhanced by QCD,

$$H_w^{\text{eff}}(\Delta S = 1) = \sqrt{2} G_F V_{ud}^* V_{us} \sum c_i O_i, \tag{3}$$

where G_F is the Fermi constant and $V_{ud}^* V_{us}$ are the Cabbibo-Kobayashi-Maskawa matrix elements. The O_i are the four-quark operators and the c_i factors are the QCD short distance Wilson coefficients [2,9]: $c_1 = -1.90 - 0.61\zeta$, $c_2 = 0.14 + 0.020\zeta$, $c_3 = c_4/5$, $c_4 = 0.49 + 0.005\zeta$, with $\zeta = V_{td}^* V_{ts}/V_{ud}^* V_{us}$. For the purpose of this work we neglect the so-called Penguin operators since their contributions are proven to be small [9]. Without QCD corrections, the Wilson coefficients would have the following values: $c_1 = -1$, $c_2 = 1/5$, $c_3 = 2/15$, $c_4 = 2/3$. In this paper we consider both possibilities and compare the resulting rates.

The techniques used to describe nonleptonic Ω^- decays (in this work we have only the $3/2^+ \to 3/2^+ + 0^-$ reaction) are known as a modified current-algebra (CA) approach. The general form of the decay amplitude reads:

$$\langle \pi(q) B'(p') | H_w^{\text{eff}} | B(p) \rangle = \overline{\mathcal{W}}_\mu(p') [(\mathcal{A} + \gamma_5 \mathcal{B}) g^{\mu\nu} + (\mathcal{C} + \gamma_5 \mathcal{D}) q^\mu q^\nu] \mathcal{W}_\nu(p). \tag{4}$$

The $\mathcal{W}(p)$ denotes the Rarita-Schwinger spinor. The parity-violating amplitudes $\mathcal{A}$ correspond to the s-wave and parity-conserving amplitudes $\mathcal{B}$ correspond to the p-wave Ω^- decays, respectively. Since the decay of $\Xi^{*0}\pi^-$ is strongly limited by phase space (momentum transfer $\simeq$ 1 MeV at the peak value of the Ξ^* mass), we will neglect the amplitudes C and $\mathcal{D}$ (d- and f- waves). The decay probability $\Gamma[\Omega^-_{3/2^+}(p) \to \Xi^{*0}_{3/2^+}(p') + \pi^-_{0^-}(q)]$ is:

$$\Gamma = \frac{|\mathbf{q}| m_f}{18\pi m_\Omega} \left\{ \left(\frac{E'}{m_f} + 1\right)\left[\frac{E'^2}{m_f^2} + \frac{E'}{m_f} + \frac{5}{2}\right]|\mathcal{A}|^2 \right. $$
$$\left. + \left(\frac{E'}{m_f} - 1\right)\left[\frac{E'^2}{m_f^2} - \frac{E'}{m_f} + \frac{5}{2}\right]|\mathcal{B}|^2 \right\}, \tag{5}$$

$$|\mathbf{q}|^2 = [(m_\Omega^2 - m_f^2 + m_\phi^2)/2m_\Omega]^2 - m_\phi^2, \tag{6}$$

$$E' = (m_\Omega^2 + m_f^2 - m_\phi^2)/2m_\Omega. \qquad (7)$$

Here Ω^- is at rest, m_f denotes the final baryon mass, and m_ϕ is the mass of the emitted meson.

We calculate the decay amplitudes by using the so-called tree-diagram approximation at the particle level, i.e., factorizable and pole diagrams plus the commutator term. The amplitudes (4) receive contributions from the commutator term and the pole diagram, respectively:

$$A_{\rm Comm} = \frac{-1}{\sqrt{2}f_\pi} a_{\Xi^*\Omega}, \qquad (8)$$

$$B_{\mathcal{P}} = \frac{-g_{\Xi^{*0}\Xi^{*-}\pi^-}}{m_\Omega - m_{\Xi^*}} a_{\Xi^*\Omega}. \qquad (9)$$

However, another state rather strongly coupled to $\Xi^*\pi$ and with a mass close to the Ω^- exists, the $\overline{\Xi}(1820, J^P = 3/2^-)$ resonance [10], whose mass fits nicely to the Gell-Man–Okubo formula for an octet of $3/2^-$ baryons. Therefore, there is a pole term in the s-wave amplitudes, $A_{\mathcal{P}}$:

$$A_{\mathcal{P}} = \frac{g_{\Xi^*\Xi^{**}\pi}}{m_\Omega - m_{\Xi^{**}}} b_{\Xi^{**}\Omega}. \qquad (10)$$

The commutator and the baryon-pole amplitudes contain weak matrix elements defined as

$$\begin{pmatrix} a_{\Xi^*\Omega} \\ b_{\Xi^{**}\Omega} \end{pmatrix} = \sqrt{2}G_F V_{ud}^* V_{us}\begin{pmatrix} \langle \Xi^*|c_i O_i^{PC}|\Omega\rangle \\ \langle \Xi^{**}|c_i O_i^{PV}|\Omega\rangle \end{pmatrix}, \qquad (11)$$

where the important parts are the Wilson coefficients and the four-quark operator matrix elements.

The factorizable contributions to s- and p-waves, are calculated by inserting the vacuum states; it is therefore a factorized product of two current matrix elements, where the decuplet-decuplet matrix element of the vector and the axial-vector currents reads:

$$\langle \Xi^*(p')|\begin{pmatrix} V \\ A \end{pmatrix}^\mu|\Omega^-(p)\rangle = g^{\Xi^*\Omega}_{\begin{pmatrix} V \\ A \end{pmatrix}} \overline{\mathcal{W}}_\nu(p)\gamma^\mu\begin{pmatrix} 1 \\ \gamma_5 \end{pmatrix}\mathcal{W}^\nu(p').$$
$$(12)$$

Summing over all factorizable contributions gives the following expressions for the amplitudes:

$$\begin{pmatrix} A \\ B \end{pmatrix}_S = \frac{G_F}{\sqrt{3}} V_{ud}^* V_{us}(m_\Omega \mp m_{\Xi^*})f_\pi g^{\Xi^*\Omega}_{\begin{pmatrix} V \\ A \end{pmatrix}}$$
$$\times[c_1 - 2(c_2 + c_3 + c_4)], \qquad (13)$$

where the $g^{\Xi^*\Omega}_{A(V)}$ represents the form-factor of the spatial component of the axial-vector (vector) current.

The total theoretical amplitudes are:

$$\mathcal{A}_{\rm th}(m_\pi^2) = A_{\rm Comm}(0) + A_{\mathcal{P}}(m_\pi^2) + A_S(m_\pi^2), \qquad (14)$$

$$\mathcal{B}_{\rm th}(m_\pi^2) = B_{\mathcal{P}}(m_\pi^2) + B_S(m_\pi^2), \qquad (15)$$

were the relative signs between commutator, pole and factorizable contributions are determined via SU(3) and the generalized Goldberger-Treiman relation.

In order to estimate the four-quark operator matrix elements entering (11), we take the Skyrme model where baryons emerge as soliton configurations of the field U of pseudoscalar mesons [11–16]. The SU(3) extended Skyrme model action is $\mathcal{L} = \mathcal{L}_\sigma + \mathcal{L}_{Sk} + \mathcal{L}_{SB} + \mathcal{L}_{WZ}$, where $\mathcal{L}_\sigma$, $\mathcal{L}_{Sk}$, $\mathcal{L}_{SB}$, and $\mathcal{L}_{WZ}$ denote the σ-model, Skyrme, symmetry breaking (SB), and Wess-Zumino (WZ) terms, respectively [16–18]. Extension of the model to the strange sector [13,15,16] is done by an isospin embedding of the static *hedgehog* ansatz into an SU(3) matrix, which is a subject of a time dependent rotation $U(\mathbf{r}, t) = A(t)\mathcal{U}(\mathbf{r})A^\dagger(t)$ by a collective coordinate matrix $A(t) \in$ SU(3). The generalized velocities are defined by $A^\dagger(t)\dot{A}(t) = \frac{i}{2}\sum_{\alpha=1}^8 \lambda_\alpha \dot{a}^\alpha$ and the profile function is interpreted as a chiral angle that parametrizes the soliton. The collective coordinates a^α are canonically quantized to generate the states that possess the quantum numbers of the physical strange baryons. In order to account for a nonzero strange quark mass the appropriate symmetry breaking terms should be included.

In this work we will use the SU(3) extended action $\mathcal{L}$ with the following set of parameters [18], $\hat{x} = 36.97$, $\beta' = -28.6$ MeV^2, $\delta' = 4.12 \times 10^7$ MeV^4, determined from the masses and decay constants of the pseudoscalar mesons. The $\hat{x}$ term is responsible for the baryon mass splittings and the admixture of higher representations in the baryon wave functions. However, we use the SU(3) symmetric baryon wave functions in the spirit of the perturbative approach to SB. Indeed, we have shown in [3,4], that the WZ and SB contributions through the weak operator matrix elements are small and introduce few % uncertainty of the decay amplitudes dominated by the Skyrme term which scales like $1/e$.

Since the coupling f_π is equal to its experimental value, for the evaluation of rare nonleptonic Ω^- decay rate the only remaining free parameter, as in [3,4], is the Skyrme charge e. It has been shown in Figure 12 of [18] that for $4.0 \leq e \leq 5.0$ the mass spectrums of **8**, **10** and $\overline{\mathbf{10}}$-plets are reasonably well described. This is the reason we are using that particular range of e further on.

The four-quark operator matrix element contribution to the commutator (8) and to the baryon-pole term (9), evaluated as a function of the Skyrme charge e [4], in units of $[10^{-8}$ GeV], is

$$a_{\Xi^{*-}\Omega^-} = \begin{cases} \text{QCD off} & \text{QCD on} & e \\ -1.86 & -3.52 & 4.00 \\ -1.31 & -2.75 & 4.75 \\ -1.22 & -2.55 & 5.00 \end{cases} \qquad (16)$$

The size of the strong coupling from pole term (9), determined via generalized Goldberger-Treiman relation,

$$g_{\Xi^{*-}\Xi^{*0}\pi^-} = \frac{m_{\Xi^{*-}} + m_{\Xi^{*0}}}{2f_\pi} g_A^{\Xi^{*-}\Xi^{0*}} \cong 9.11 g_A^{pn}, \quad (17)$$

is close to the estimate given in [19].

The strong coupling $g_{\Xi^{**}\Xi^*\pi}$, which enters in the calculation of $A_{\mathcal{P}}$, we extract from experimental value for partial width $\Gamma(\Xi^{**} \to \Xi^*\pi)$. Using strong effective Lagrangian for spin $3/2^-$ baryon Ξ^{**}

$$\mathcal{L}_{int}(\Xi^{**} \to \Xi^*\pi) = g_{\Xi^{**}\Xi^*\pi} \overline{W}_\mu(p') \tilde{W}^\mu(p) \phi_\pi(q), \quad (18)$$

where $\tilde{W}^\mu$ denotes the $3/2^-$ Rarita-Schwinger spinor, we obtain for Ξ^{**} at rest,

$$\Gamma(\Xi^{**} \to \Xi^*\pi) = \frac{g_{\Xi^{**}\Xi^*\pi}^2 |\mathbf{q}|}{36\pi m_{\Xi^*}^2 m_{\Xi^{**}}} (E' + m_{\Xi^*})$$
$$\times [2E'^2 + 2E'm_{\Xi^*} + 5m_{\Xi^*}^2]. \quad (19)$$

The E' and $|\mathbf{q}|$ are given in (6) and (7). From experiment [6] we have $\Gamma_{tot}(\Xi^{**}) = (24 \pm 6)$ MeV and

$$\Gamma(\Xi^{**} \to \Xi^*\pi)/\Gamma_{tot}(\Xi^{**}) = (0.30 \pm 0.15). \quad (20)$$

The generalized Goldberger-Treiman relation applied to (10) gives the product $b_{\Xi^{**}\Omega} g_{\Xi^{**}\Xi^*}^A$ which does not depend on the phase of the exchanged Ξ^{**} in the pole diagram. It has been found in [20] that the sign of an axial-vector matrix element, between particles of the opposite parity, is negative. Then from (19) and (20) we find $g_{\Xi^{**}\Xi^*\pi} = -(0.48 \pm 0.14)$, a small negative value with large uncertainty.

The weak matrix elements of the same spin but opposite parity baryons are found to be about the same as the ones including ground states only [20], i.e., $b_{\Xi^{**}\Omega} \cong a_{\Xi^*\Omega}$. This together with (8) and (10) gives

$$A_{\mathcal{P}} = -(0.44 \pm 0.12)A_{Comm}, \quad (21)$$

up to the sign, similar to the conclusion in [19]. Since the resonance Ξ^{**} is known with quite a large error, the same error is also present in (21).

We proceed with the computation of the vector and axial-vector current form-factors $g_{V(A)}^{\Xi^*\Omega}$, in terms of $g_{V(A)}^{pn}$, via SU(3) Clebsh-Gordan coefficients. Using from experiment, $g_A^{pn}/g_V^{pn} = 1.26$, we find an agreement with the SU(6) result [19], as we should:

$$g\binom{\Xi^{*-}\Omega^-}{\begin{array}{c}V\\A\end{array}} = g\binom{\Xi^{*0}\Omega^-}{\begin{array}{c}V\\A\end{array}} \cong \binom{1.1}{0.7}. \quad (22)$$

Factorizable amplitudes A_S and B_S are very slowly changing functions, via the g_A^{pn}/g_V^{pn}, with respect to the Skyrme charge e, (for details see Figs. 2–4 in [18]). For QCD off/on we found from (13) and (22):

$$A_S^{off} = -5.65 \times 10^{-8}, \qquad A_S^{on} = -6.32 \times 10^{-8},$$
$$B_S^{off} = -90.1 \times 10^{-8}, \qquad B_S^{on} = -100.8 \times 10^{-8}. \quad (23)$$

The mean values of theoretical amplitudes A_{Comm}, $A_{\mathcal{P}}$ and $B_{\mathcal{P}}$ for QCD off/on, as functions of the Skyrme charge e, are presented in Table I.

Rare nonleptonic two-body Ω^- decay partial width formulas

$$\Gamma_{th}(\Omega^- \to \Xi^{*0}\pi^-)$$
$$= \{9.000617|A_S + A_{Comm} + A_{\mathcal{P}}|^2$$
$$+ 0.000148|B_S + B_{\mathcal{P}}|^2\}2.697 \times 10^{-4} \text{ GeV}, \quad (24)$$

and Eqs. (8)–(23), in units of $[10^{-18}$ GeV], gives

$$\Gamma_{th}(\Omega^- \to \Xi^{*0}\pi^-) = \begin{cases} \text{QCD off} & \text{QCD on} & e \\ 1.19^{+2.65}_{-1.09} & 17.81^{+16.57}_{-10.44} & 4.00 \\ 0.01^{+0.38}_{-0.01} & 6.84^{+8.36}_{-4.77} & 4.75 \\ 0.06^{+0.54}_{-0.06} & 4.83^{+6.65}_{-3.60} & 5.00 \end{cases} \quad (25)$$

which, compared with the experiment [6],

$$\Gamma_{exp}(\Omega^- \to \Xi^{*0}\pi^-) = \left(5.1 \begin{array}{c} +4.1 \\ -1.6 \end{array}\right) \times 10^{-18} \text{ GeV}, \quad (26)$$

shows the following:

(a) In this dynamical scheme framework, contrary to the nonleptonic hyperon and Ω^- decays [3,4], the factorizable contributions turn out to be very important for the rare $\Omega^- \to \Xi^{*0}\pi^-$ nonleptonic decay. The opposite signs between the commutator and pole term (21) and the factorizable contributions (23) becomes an essential feature of our dynamical scheme, leading to the internal cancellation within the $\mathcal{A}$ and $\mathcal{B}$ amplitudes and bringing theoretical estimate (25) closer to the experiment.

(b) From theoretical rate (25) it is clear that in this dynamical scheme the QCD enhancement is crucial.

(c) An important feature of the particular rare decay mode $\Omega^- \to \Xi^{*0}\pi^-$ lies in the fact that it is strongly limited by phase space and that the decay transition occurs at threshold. Thus, π^- in the final state is almost at rest.

(d) First consequence of (c) is the emitted pion is really soft and the so-called soft pion limit theorem is very well satisfied. This is the reason for the commutator term

TABLE I. The amplitudes A_{Comm}, $A_{\mathcal{P}}$, $B_{\mathcal{P}}$ in units $[10^{-8}]$ contributing to the rare nonleptonic two-body Ω^- decay rate.

	QCD off			QCD on		
e	A_{Comm}	$A_{\mathcal{P}}$	$B_{\mathcal{P}}$	A_{Comm}	$A_{\mathcal{P}}$	$B_{\mathcal{P}}$
4.00	14.14	−6.22	155.6	26.76	−11.77	294.5
4.75	9.96	−4.38	109.6	20.91	−9.20	230.1
5.00	9.28	−4.08	102.1	19.39	−8.53	213.4

dominance in the s-wave amplitude $\mathcal{A}$. Second consequence of (c) is an overall phase space enhancement, by factor of 6.1×10^4, for the s-waves. Owing to this, although the p-wave amplitudes are more than 1 order of magnitude larger than the corresponding s-waves, the total contribution of $\mathcal{B}$ amplitudes to the rate (25) is very small.

(e) Inspection of the result (25) shows the importance of the e dependence and that $e = 4.75$ with QCD corrections switched on gives the central fit for the rare $\Omega^- \to \Xi^{*0}\pi^-$ nonleptonic decay. Applying the same e and QCD switched on to the Ω^- nonleptonic decays from [4], we first find $a_{\Lambda\Xi^0} = 4.85 \times 10^{-8}$ GeV. This together with the matrix element $a_{\Xi^-\Omega^-} = -2.75 \times 10^{-8}$ GeV and factorizable contributions from Table 1 in [4] produces the Ω^- nonleptonic decay amplitudes: $\mathcal{B}_{\text{th}}(\Omega_K) = 6.02$, $\mathcal{B}_{\text{th}}(\Omega_-) = 1.48$ and $\mathcal{B}_{\text{th}}(\Omega_0^-) = 1.29$, in units of $[10^{-6} \text{ GeV}^{-1}]$, which are closer to the conclusions in [21] and in better agreement with experiment than the ones obtained in [4]. The same should hold for the p-wave nonleptonic hyperon decays [3].

Although our dynamical scheme is the same as the one in [19], the situation has changed during the past 20 years. The important improvements happened in the areas of QCD corrections to the effective weak Hamiltonian and in the further development of the Skyrme model. The latter becomes a good candidate for the correct description of the higher SU(3) representation multiplets [18,22]

and a good tool for evaluation of the nonperturbative quantities like the four-quark operator matrix elements between different baryon states. Comparing the result (25) with Table 1 from [19], we conclude that the dynamical scheme is working well in both cases. In this work we are using well accepted results for the next to leading order computations of the QCD corrections [9]. According to [9] the errors generated by different choice of renormalization scheme and scale are less than 10%. Overall QCD corrections are more stable and smaller today; penguins are about 1 order of magnitude smaller and also $c_1 \simeq -3 \to -1.9$ [9]. The Skyrme model provides us with four-quark operator matrix elements, in a different way than the MIT bag model did in [1,19], producing rough agreement with measurement (26). Finally, according to (e), as a bonus we obtain better agreement with experiment for the nonleptonic Ω^- and the p-wave hyperon decays.

Obviously, not all details are under full control. For example, m_s corrections are neglected and the uncertainties contained in the input parameters (20) are unfortunately large. Nevertheless, the QCD corrected weak Hamiltonian H_w^{eff}, supplemented by the Skyrme model, for the mean values of the input parameters, leads to the correct description of the rare $\Omega^- \to \Xi^0(1530)\pi^-$ decay.

[1] H. Galić *et al.*, Phys. Lett. **89B**, 249 (1980); D. Tadić and J. Trampetić, Phys. Rev. D **23**, 144 (1981); **30**, 1990 (1984).

[2] J. F. Donoghue, H. Golowich, and B. R. Holstein, *Dynamics of the Standard Model* (Cambridge Univ. Press, Cambridge, England, 1994).

[3] G. Duplancic *et al.*, Phys. Rev. D **64**, 097502 (2001).

[4] G. Duplancic *et al.*, Phys. Rev. D **65**, 054001 (2002).

[5] In the further we should note $\Xi^0_{3/2^+}(1530) \equiv \Xi^*$.

[6] D. E. Groom *et al.*, Eur. Phys. J. C **15**, 1 (2000); N. Solomey, Nucl. phys., Sect. B Proc. suppl. **115**, 54 (2003).

[7] J. Trampetić, Phys. Lett. **144B**, 250 (1984); M. Praszałowicz and J. Trampetić, Phys. Lett. **161B**, 169 (1985); Fizika (Zegreb) **18**, 391 (1986).

[8] D. I. Diakonov *et al.*, Nucl. Phys. **B323**, 53 (1989).

[9] A. J. Buras, hep-ph/9806471.

[10] In the further we should note $\Xi^-_{3/2^-}(1820) \equiv \Xi^{**}$.

[11] T. H. R. Skyrme, Proc. R. Soc. London A **260**, 127 (1961).

[12] G. Adkins *et al.*, Nucl. Phys. **B228**, 552 (1983).

[13] E. Guadagnini, Nucl. Phys. **B236**, 35 (1984).

[14] H. Yabu and K. Ando, Nucl. Phys. **B301**, 601 (1988).

[15] P. O. Mazur *et al.*, Phys. Lett. **147B**, 137 (1984); S. Jain and S. Wadia, Nucl. Phys. **B258**, 713 (1985).

[16] H. Weigel, Int. J. Mod. Phys. A **11**, 2419 (1996).

[17] N. N. Scoccola, Phys. Lett. B **428**, 8 (1998); D. Dumm *et al.*, Phys. Rev. D **62**, 014001 (2000).

[18] G. Duplancic, H. Pasagic, and J. Trampetic, J. High Energy Phys. 07 (2004) 027.

[19] M. Lusignoli and A. Pugliese, Phys. Lett. **132B**, 178 (1983).

[20] S. Pakvasa and J. Trampetić, Phys. Lett. **126B**, 122 (1983); M. Bonvin, Nucl. Phys. **B38**, 241 (1984).

[21] G. Eilam *et al.*, Phys. Rev. D **53**, 3629 (1996); B. Borasoy and B. R. Holstein, Phys. Rev. D **60**, 054021 (1999).

[22] V. Kopeliovich, Phys. Lett. B **259**, 234 (1991); H. Weigel, Eur. Phys. J. A **2**, 391 (1998); H. Walliser and V. B. Kopeliovich, J. Exp. Theor. Phys. **97**, 433 (2003); V. Kopeliovich, hep-ph/0310071; M. Praszalowicz, Phys. Lett. B **575**, 234 (2003); N. Itzhaki *et al.*, Nucl. Phys. **B84**, 264 (2004); B. Wu and B. Q. Ma, Phys. Lett. B **586**, 62 (2004); J. Ellis *et al.*, J. High Energy Phys. 05 (2004) 002; G. Duplancic and J. Trampetic, Phys. Rev. D **69**, 117501 (2004); H. Weigel, Eur. Phys. J. A **21**, 133 (2004); G. Duplancic, H. Pasagic, and J. Trampetic, hep-ph/0404193.