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INTRODUCTION TO CHIRAL DOUBLING OF HEAVY–LIGHT HADRONS*

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In this lectures we review the main idea of the chiral doublers scenario, originating from simultaneous constraints of chiral symmetry and of heavy quark spin symmetry on effective theories of heavy–light hadrons. We discuss chiral doublers for mesons and (briefly) chiral doublers for baryons in light of recent experimental data.

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1. Introduction

These notes are organized as follows: first, we outline recent experimental data, which triggered the renewed interest in physics of heavy hadrons (Section 2). Then, after briefly mentioning the plethora of theoretical proposals to describe these states, we choose the particular one (in our opinion, most appealing) — the chiral doublers interpretation for the new states. We show, how one can "guess" the unique leading order effective Lagrangian incorporating both symmetries of the heavy and light quarks (Section 3). Then we present general arguments, why the phenomenon of chiral doubling seems to be the generic pattern of the QCD. Finally we propose the classification of the observed to-date heavy–light mesons in the chiral doubling scheme. In the last part (Section 4) we briefly discuss the heavy–light baryons. Despite the chiral scenario for baryons seems to be quite appealing, the present lack of solid experimental data for heavy–light baryons makes the verification difficult. In the context of the chiral doubling, we speculate also on recent exotic baryonic signal from H1 experiment, interpreted as a signal for charmed pentaquark.

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2. New experimental results on open charm — new sensations

Recently, experimental physics of hadrons with open charm has provided several spectacular discoveries which surprised almost everyone. Let us briefly remind these experiments:

- In April 2003 BaBar Collaboration [1] has announced new, narrow meson $D^{\star}_{sJ}(2317)^{+}$, decaying into D^{+}_{s} and π^{0}. In May 2003 this observation was confirmed by CLEO Collaboration [2], which also noticed another narrow state, $D_{sJ}(2460)^{+}$, decaying into $D^{\star}_{s}$ and π^{0}. Both states were confirmed by Belle Collaboration [3], and finally, the CLEO observation was also confirmed by BaBar [4].

- In July 2003 Belle Collaboration measured the narrow excited states D_1, D_2 with foreseen quantum numbers $(1^{+}, 2^{+})$, and provided the first evidence for two new, broad states $D^{\star}_{0}$ $(2308 \pm 17 \pm 15 \pm 28)$ and $D^{'}_{1}$ $(2427 \pm 26 \pm 20 \pm 17)$ [5]. Both of them are approximately 400 MeV above the usual $D_0, D^{\star}$ states and seem to have opposite to them parity. Similar observation of the $D^{\star}_{0}$ was observed by FOCUS Collaboration [6] and also CLEO has noticed second state $D^{'}_{1}$ [2]. Above states were seen in decays mode with pion, *i.e.* $D^{\star 0}_{0} \rightarrow D^{+}\,\pi^{-}$, $D^{\star +}_{0} \rightarrow D^{0}\,\pi^{+}$ and ${D^{'}_{1}}^{0} \rightarrow D^{\star +}\pi^{-}$.

- Selex Collaboration announced very recently a new, surprisingly narrow state $D^{+}_{sJ}(2632)$ [7], which mainly appeared in $D^{+}_{sJ} \rightarrow D^{+}_{s}\eta$ decays.

- H1 experiment at DESY has announced [8] a signature for charmed pentaquark Θ^{0}_{c} $(\bar{c}udud)$ at mass 3099 MeV, *i.e.* approximately 400 MeV higher than the expected estimates known in the literature [9–12].

Till today last two states were neither confirmed nor falsified by other experiments.

3. Mesons

We visualize the schematic spectroscopy of new mesons on Fig. 1. The above states and in particular the decay patterns of all these particles challenged standard estimations based on quark potential models (QM) and triggered a renewal of interest on charmed hadrons spectroscopy among several theorists. Why these states are so surprising? If we consider an infinitely heavy quark, its spin decouples from the angular momentum of the light object j_l. For the lowest partial wave of light quark $l = 0$, j_l equals to the spin of the light quark. By adding light and heavy components, we get a pair of states $0^{-}, 1^{-}$. Next partial wave corresponds to $l = 1$, therefore we have two possibilities $j_l = 3/2$ and $j_l = 1/2$. Adding the spin of the heavy quark to j_l we get two pairs $1^+, 2^+$ and $0^+, 1^+$, respectively. The observations of BaBar, CLEO and Belle point that new $D_{\rm s}$ states (2317) and (2460) match the spin-parity pattern of the last pair. The first puzzle of these states was therefore not their presence, but the value of their masses. QM predictions were placing these states ca 150 MeV higher than observed, *i.e.* above corresponding mass thresholds for DK and D^*K. Such states were therefore expected to be broad. The new states were however surprisingly narrow with width below 10 MeV. Second puzzle was the pattern of splitting between the particles of opposite parity: the mass split between 0^+ and 0^- turned out to be almost identical to the mass split between 1^+ and 1^-. The third challenge was to understand the decays — hadronic (basically one pion emissions) and electromagnetic.

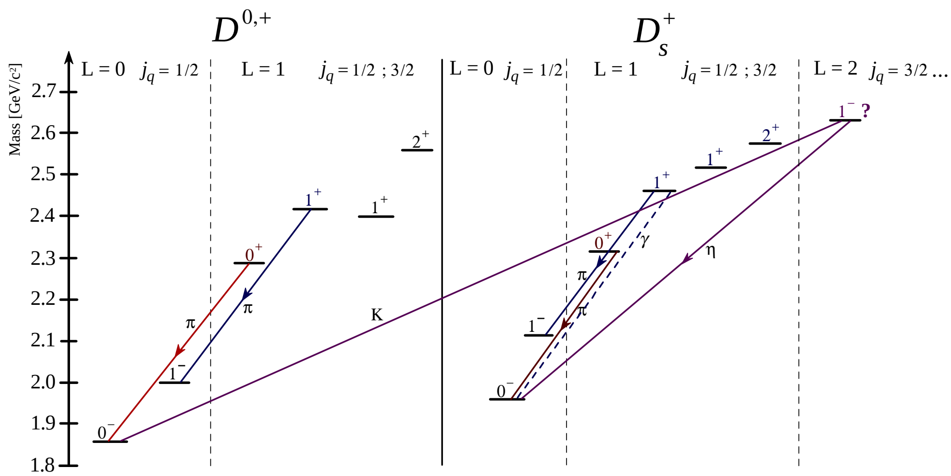

Fig. 1. Spectroscopy of $D_0^\star$,$D_1^{'}$ (left side) and $D_{sJ}^\star$,D_{sJ},D_{sJ}^+ mesons (right side). The lines show observed decays for new particles.

These three challenges triggered an interest on charmed hadrons spectroscopy among several theorists [13, 14]. In the above mentioned works, new states were interpreted either as tetraquarks, or as "molecular configurations" alike $D\pi$ atoms or DK molecules, or as resonant states forced by unitarization and chiral symmetry. There were also works trying to interpret these states in the framework of modified QM or verifying their properties via lattice simulation. Last but not least, a decade old idea [15,16] of chiral doubling was brought to attention [17].

3.1. Chiral (χ) doublers scenario

With respect to $\Lambda_{\rm QCD}$, the fundamental scale of Quantum Chromodynamics, strong interactions involve three light flavors ($q = u, d, s$) and three heavy flavors ($Q = c, b, t$), (see Fig. 2). It is instructive to consider the limits:

- $m_q \to 0$
- $m_Q \to \infty$.

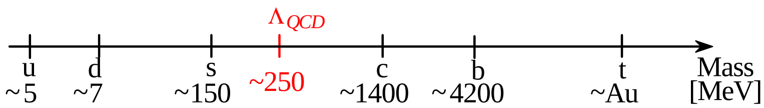

Fig. 2. Schematic QCD mass scales.

Both limits (massless quarks and infinitely heavy mass) unravel essential symmetries of the interactions. The light sector (massless light quark limit) is characterized by the spontaneous breaking the chiral symmetry (SBχS). Vacuum state is respecting only vector part of the symmetry, *i.e.* $\mathrm{SU_V}(N_q) \times \mathrm{SU_A}(N_q) \to \mathrm{SU_V}(N_q)$, whereas axial symmetry is broken, as a result we have massless Goldstone's excitations for each broken generator.

The heavy sector (infinite heavy quark mass limit) exhibits heavy quark symmetry (Isgur–Wise symmetry) [18]. In this limit, dynamics of the heavy quark becomes independent of its spin and mass. As a result of such a limit the masses of the pseudoscalar (0^-) and vector (1^-) mesonic states, including heavy quark become degenerate.

heavy–light mesons are the simplest objects subjected to the simultaneous restrictions of both above-mentioned symmetries. Constraints from both symmetries enforce the form of the effective interaction of such mesons. An explicit answer from theoretical point of view was found in 1992 and 1993 [15, 16] and the major consequence of derivation proposed was that the interaction requires an introduction of *chiral partners*. Below we present the argument, how one can guess/derive such an interaction using the approximate bosonization scheme for QCD.

3.2. Schematic constructions for heavy–light mesons $\bar{q}Q$

For the simplicity of the algebra we restrict our discussion to two light flavors $q = (u, d)$ and a heavy flavor ($Q = c$). The generalization to ($Q = b$) and ($q = s$) is straightforward. If the mass of the heavy quark is infinitely large, then the heavy quark momentum is large and conserved $P_\mu = m_Q v_\mu$. In this limit, we have a velocity superselection rule [19], *i.e.* we have a different heavy quark (antiquark) field $Q_v^{\pm}(x)$ for each velocity v. To display this we follow Georgi [19] and define heavy quark field as

$$Q(x) = \frac{1+\not{v}}{2} e^{-im_Q v\cdot x} Q_v^{+}(x) + \frac{1-\not{v}}{2} e^{im_Q v\cdot x} Q_v^{-}(x)\,, \tag{1}$$

where $(1 \pm \not{v})/2$ are projection operators originating from standard $(\not{p} \pm m_Q)/2m_Q$. As a result, the free QCD action in terms of light and heavy quark fields reads

$$S = \sum_v \int d^4x \ \left(\bar{q}(i\not{\partial} - m_q)q + \bar{Q}_v(i\not{v}\, v\cdot\partial)Q_v\right)\,. \tag{2}$$

Notice that our action (2) is flavor $\mathrm{U(2)_L \times U(2)_R}$ symmetric (for $m=0$) and invariant under independent spin rotations of the quark and the antiquark. After applying approximate bosonization schemes [20] to the heavy–light system we can generate an effective action as a gradient expansion in the slowly varying fields that intermingles heavy–light dynamics [15], alike similar schemes lead to the effective mesonic Lagrangians (sigma models) for light flavors. We denote the heavy meson fields as

$$\hat{H}_{\pm} = \frac{1+\not{v}}{2}(\gamma^{\mu}\hat{P}^{*}_{\mu,\pm} + i\gamma_5\hat{P}_{\pm})\gamma_5^{\pm} + (\mathrm{h.c.})\,, \tag{3}$$

where

$$\begin{aligned} \hat{P}^{a}_{+} &\sim \bar{q}^{a}_{\mathrm{R}}\, Q_v\,, \\ \hat{P}^{*a}_{\mu,+} &\sim \bar{q}^{a}_{\mathrm{R}}\, \gamma_\mu\, Q_v \end{aligned} \tag{4}$$

are the bare pseudoscalar and bare vector heavy mesons with light chirality and v_μ is the velocity of the heavy quark ($v^2=1$). Note that changing $(+ \to -)$ corresponds to $(\mathrm{R} \to \mathrm{L})$.

The action Eq. (2) can be written in the form

$$\begin{aligned} S \;=\; \sum_v \int \bar{\psi}\Big(&\mathbf{1}_2 i\not{\partial} + \mathbf{1}_3 i\not{v} v\cdot\partial + \mathbf{1}_2(\hat{\not{L}}\gamma_5^{+} + \hat{\not{R}}\gamma_5^{-}) - \mathbf{1}_2(M\gamma_5^{+} + M^{\dagger}\gamma_5^{-}) \\ &+\hat{H}_{+} + \hat{H}_{-}\Big)\psi \end{aligned} \tag{5}$$

with quark field $\psi = (q; Q_v)$ and bare light vector fields

$$\hat{L}_\mu \sim \bar{q}_{\mathrm{L}}\gamma_\mu q_{\mathrm{L}}\,, \qquad \hat{R}_\mu \sim \bar{q}_{\mathrm{R}}\gamma_\mu q_{\mathrm{R}} \tag{6}$$

valued in $\mathrm{U(2)_L}$ and $\mathrm{U(2)_R}$, respectively.

In above expression we used also the projectors onto the light and heavy sectors in the form $\mathbf{1}_2 = \mathrm{diag}(1,1,0)$, $\mathbf{1}_3 = \mathrm{diag}(0,0,1)$ and the short-hand notation $\gamma_5^{\pm} \equiv \frac{1}{2}(1 \pm \gamma_5)$. The $\hat{P}$'s in Eq. (4) are off diagonal in flavor space, so that $H = H_a T^a$ with $a = 1,2$ and $T^1 = (\lambda_4 - i\lambda_5)/2$, $T^2 = (\lambda_6 - i\lambda_7)/2$. Note that λ are standard Gell-Mann matrices. The effect of spontaneous breakdown of the chiral symmetry is introduced via matrix M, chosen as

$$M \;=\; \xi_{\mathrm{L}}^{\dagger}\; \Sigma\; \xi_{\mathrm{R}} \tag{7}$$

with

$$\xi_{\mathrm{L}}^{\dagger} = \xi_{\mathrm{R}} = \exp\left(\frac{i\vec{\pi}\vec{\tau}}{2f_{\pi}}\right) , \tag{8}$$

where $f_{\pi} = 93$ MeV is the pion decay constant and $\vec{\tau}$ are standard Pauli matrices. In the vacuum Σ is diagonal and constant along the light directions (u,d). If we define the constituent (dressed) quark field χ by the relation to the bare quark field ψ,

$$\begin{aligned} \chi_{\mathrm{L}} &= (\xi_{\mathrm{L}} q_{\mathrm{L}}; Q_v)\,, \\ \chi_{\mathrm{R}} &= (\xi_{\mathrm{R}} q_{\mathrm{R}}; Q_v)\,, \end{aligned} \tag{9}$$

we can rewrite the "bosonized" action given by Eq. (5) in terms of field χ. Then the action reads

$$\begin{aligned} S \;=\; \sum_v \int d^4x \overline{\chi}\, [\, & \mathbf{1}_2 i\partial\!\!\!/ + \mathbf{1}_3 i\rlap{/}v v\cdot\partial + \mathbf{1}_2(\xi_{\mathrm{R}} i\partial\!\!\!/ \xi_{\mathrm{R}}^{\dagger}\gamma_5^{+} + \xi_{\mathrm{L}} i\partial\!\!\!/ \xi_{\mathrm{L}}^{\dagger}\gamma_5^{-}) \\ & + \mathbf{1}_2(\xi_{\mathrm{L}}\hat{L}\!\!\!/ \xi_{\mathrm{L}}^{\dagger}\gamma_5^{+} + \xi_{\mathrm{R}}\hat{R}\!\!\!/ \xi_{\mathrm{R}}^{\dagger}\gamma_5^{-}) - \mathbf{1}_2(\Sigma\gamma_5^{+} + \Sigma^{+}\gamma_5^{-}) \\ & + \overline{H} + H + \overline{G} + G\,]\, \chi\,. \end{aligned} \tag{10}$$

An immediate feature of the above action is the mandatory appearance of another heavy meson field denoted by G, which explicit form reads

$$\begin{aligned} H &= \frac{1+\rlap{/}v}{2}(\gamma^{\mu} D_{\mu}^{*} + i\gamma_5 D) \\ &= \frac{1+\rlap{/}v}{2}\left(\gamma^{\mu}(P_{\mu,+}^{*}\xi_{\mathrm{R}}^{+} + P_{\mu,-}^{*}\xi_{\mathrm{L}}^{+}) + i\gamma_5(P_{+}\xi_{\mathrm{R}}^{+} + P_{-}\xi_{\mathrm{L}}^{+})\right) , \\ G &= \frac{1+\rlap{/}v}{2}(\gamma^{\mu}\gamma_5 \tilde{D}_{\mu}^{*} + \tilde{D}) \\ &= \frac{1+\rlap{/}v}{2}\left(\gamma^{\mu}\gamma_5(P_{\mu,+}^{*}\xi_{\mathrm{R}}^{+} - P_{\mu,-}^{*}\xi_{\mathrm{L}}^{+}) + (P_{+}\xi_{\mathrm{R}}^{+} - P_{-}\xi_{\mathrm{L}}^{+})\right) , \end{aligned} \tag{11}$$

where D and $D^{\star}$ in field H represent the pseudoscalar (0^-) and the vector (1^-) mesons fields, respectively, which annihilate the $s_{\ell} = \frac{1}{2}$ meson multiplet. Field G refers to new $(\tilde{D}, \tilde{D}^{\star})$ multiplet with the constrain $v^{\mu}\tilde{D}_{\mu}^{\star} = 0$, *i.e.* to the chiral doubler with spin-parity assignment $(0^+, 1^+)$. Dressed fields L/R follows from

$$\begin{aligned} L_\mu &= \xi_\mathrm{L} \hat{L}_\mu \xi_\mathrm{L}^\dagger + i \xi_\mathrm{L} \partial_\mu \xi_\mathrm{L}^\dagger \\ R_\mu &= \xi_\mathrm{R} \hat{R}_\mu \xi_\mathrm{R}^\dagger + i \xi_\mathrm{R} \partial_\mu \xi_\mathrm{R}^\dagger \,. \end{aligned} \tag{12}$$

Under heavy quark spin symmetry $\mathrm{SU}(2)_Q$ (denoted by S)

$$\begin{aligned} H &\to SH \,, \\ G &\to SG \end{aligned} \tag{13}$$

and chiral $\mathrm{SU}(2)_\mathrm{L} \times \mathrm{SU}(2)_\mathrm{R}$ (denoted by U) transforms

$$\begin{aligned} H &\to HU^\dagger \,, \\ G &\to GU^\dagger \,. \end{aligned} \tag{14}$$

It is also convenient to introduce

$$\begin{aligned} \bar{H} &= \gamma^0 H^\dagger \gamma^0 \,, \\ \bar{G} &= \gamma^0 G^\dagger \gamma^0 \end{aligned} \tag{15}$$

and transformations for them

$$\begin{aligned} \bar{H} &\to U\bar{H} \,, \\ \bar{G} &\to U\bar{G} \,, \end{aligned} \tag{16}$$

$$\begin{aligned} \bar{H} &\to \bar{H}S^\dagger \,, \\ \bar{G} &\to \bar{G}S^\dagger \,. \end{aligned} \tag{17}$$

After substituting new dressed fields from Eq. (12) to Eq. (10) we arrive at the effective action

$$\begin{aligned} S = \sum_v \int d^4x \overline{\chi} \big(& \mathbf{1}_2 (i \not{\nabla}_\mathrm{L} - \varSigma) \gamma_5^- + \mathbf{1}_2 (i \not{\nabla}_\mathrm{R} - \varSigma) \gamma_5^+ + \mathbf{1}_3 i \not{v} v \cdot \partial \\ & + H + \overline{H} + G + \overline{G} \big) \chi \end{aligned} \tag{18}$$

where covariant derivatives are

$$\begin{aligned} \nabla_\mathrm{L} &= \partial - iL \,, \\ \nabla_\mathrm{R} &= \partial - iR \,. \end{aligned} \tag{19}$$

This is the starting point for the derivative expansion, with quark propagator of the form

$$\boldsymbol{S} = \left(\mathbf{1}_2(i\not\nabla_{\rm L} - \Sigma)\gamma_5^- + \mathbf{1}_2(i\not\nabla_{\rm R} - \Sigma)\gamma_5^+ + \mathbf{1}_3 i\not v v\cdot\partial\right)^{-1} \tag{20}$$

and generic heavy–light part of the Lagrangian density

$$\mathcal{L} = \bar{\chi}\left(\boldsymbol{S}^{-1} + H + \bar{H} + G + \bar{G}\right)\chi\,. \tag{21}$$

We integrate now the partition function over the fermions, and then we rewrite the resulting determinant using the well-known trick

$$e^{(\ln\det A)} = e^{(\mathrm{Tr}\,\ln A)}\,. \tag{22}$$

Expansion of the logarithm to the second order gives the effective action, *e.g.* for H fields we arrive at induced at one loop level action

$$\begin{aligned} S_H &= \frac{1}{2}N_c\mathrm{Tr}\,\left(\mathbf{1}_2\Delta_q H\mathbf{1}_3\Delta_Q\bar{H}\right) \\ &\quad -\frac{1}{2}N_c\mathrm{Tr}\,\left(\mathbf{1}_2\Delta_q(\not V\Delta_q + \not A\Delta_q\gamma_5)H\mathbf{1}_3\Delta_Q\bar{H}\right) + \cdots\,, \end{aligned} \tag{23}$$

where

$$\Delta_q = (i\not\partial - \Sigma)^{-1}\,, \qquad \Delta_Q = (i\not v v\cdot\partial)^{-1} \tag{24}$$

are dressed light and heavy flavor propagators, respectively and the functional trace includes tracing over space, flavor and spin indices. The ellipsis in Eq. (23) stands for higher insertions of vectors and axials. Note the light–light and heavy–light quark dynamics follows from the dressed action (18) through a derivative expansion. The resulting effective action for the H sector reconstructs (modulo chiral mass term) the original construction by [21]. Our expansion of (18) will be understood in the sense of $m_Q/\Lambda \to \infty$ ($m_Q \to \infty$). A similar action appears for the heavy chiral partners G's.

After regularizing and renormalizing the resulting one loop integrals one arrives at the final form of the effective action

$$\begin{aligned} \mathcal{L}_v^H &= -\frac{i}{2}\mathrm{Tr}\,(\bar{H}v^\mu\partial_\mu H - v^\mu\partial_\mu\bar{H}\,H) \\ &\quad +\mathrm{Tr}\,V_\mu\bar{H}Hv^\mu - g_H\mathrm{Tr}\,A_\mu\gamma^\mu\gamma_5\bar{H}H + m_H\mathrm{Tr}\,\bar{H}H\,, \end{aligned} \tag{25}$$

$$\begin{aligned} \mathcal{L}_v^G &= +\frac{i}{2}\mathrm{Tr}\,(\bar{G}v^\mu\partial_\mu G - v^\mu\partial_\mu\bar{G}\,G) \\ &\quad -\mathrm{Tr}\,V_\mu\bar{G}Gv^\mu - g_G\mathrm{Tr}\,A_\mu\gamma^\mu\gamma_5\bar{G}G + m_G\mathrm{Tr}\,\bar{G}G\,. \end{aligned} \tag{26}$$

Note that the masses m_H and m_G are of order m_Q^0.

Chiral partners communicate with each other via light axial currents

$$\mathcal{L}_v^{HG} = g_{GH}\mathrm{Tr}\,(\gamma_5 \bar{G} H \gamma^\mu A_\mu) + (\mathrm{h.c.})\,, \tag{27}$$

where g_{GH} is the coupling constant governing the $(0^+, 1^+) \to (0^-, 1^-)$ "π" transitions and we do not have vector mixing because of the parity.
The axial A_μ and vector V_μ read, respectively

$$A_\mu \;=\; \frac{i}{2}(\xi^\dagger \partial_\mu \xi - \xi \partial_\mu \xi^\dagger)\,, \tag{28}$$

$$V_\mu \;=\; \frac{1}{2}(\xi^\dagger \partial_\mu \xi + \xi \partial_\mu \xi^\dagger)\,. \tag{29}$$

The main consequences of above derivation are as follows:

- Chiral symmetry and heavy quark symmetry (IW symmetry) require the introduction of *chiral partners*. Chiral partners (G) are parity duplications for the standard multiplet (H).
- The mass splitting between the usual multiplet H and chiral partner G imply the mass relation to order $m_Q^0 N_c^0$

$$m_G - m_H = m(\tilde{D}^\star) - m(D^\star) = m(\tilde{D}) - m(D) = \mathcal{O}(\Sigma)\,, \tag{30}$$

  where Σ denotes one loop heavy meson self-energy [15–17]. Look at the diagram shown in Fig. 3.

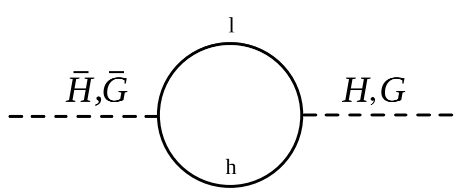

Fig. 3. One-loop contribution to 2-point $\bar{H}H$, $\bar{G}G$ functions. Here l stands for light quark and h for heavy quark.

- Technically, the difference for chiral masses originates from the γ_5 difference in the definition of the fields H and G. In other words, it is sensitive to the parity content of the heavy–light field since $H\not{v} = -H$ and $G\not{v} = +G$. The result is the mass gap between the heavy–light mesons of opposite chirality. This unusual contribution of the chiral quark mass stems from the fact that it tags to the velocity $H\not{v}\bar{H}$ of the heavy field and is therefore sensitive to parity.

- If we have restored the spontaneous breakdown of the chiral symmetry, the mass gap disappears. Therefore (in the chiral limit), such a mass gap is an order parameter for the spontaneous breakdown of the chiral symmetry, and actually this mass gap could be even used as a definition of rather elusive concept of dressing the current mass of the quark. Surprisingly, heavy–light system seems to be quite adequate for probing the chiral properties of the vacuum (since only one quark gets dressed, contrary to two quarks being dressed for light mesons).

- Such a generic phenomenon cannot be model dependent, one therefore expects low-energy theorems for chiral doublers. Indeed, this is the case. Since the leading term in the axial current comes from one-pion, $A^a_\mu \sim 1/f_\pi \partial_\mu \pi^a$, integrating by parts the expectation value of the mixed-term in the Lagrangian gives the Goldberger–Treiman-type relation between the doublers, *i.e.*

$$m_G - m_H \sim g_{HG} f_\pi \,. \tag{31}$$

  This means that a *small scale* of order of 100 MeV (originating from pion decay constant, *i.e.* the quantity related directly to the properties of the QCD vacuum) appears naturally even for very heavy mesons (*e.g.* for B mesons).

3.3. D-cubes

In this subsection we visualize the consequences of the chiral doublers scenario for mesons in the form of the cartoon, see Fig. 4.

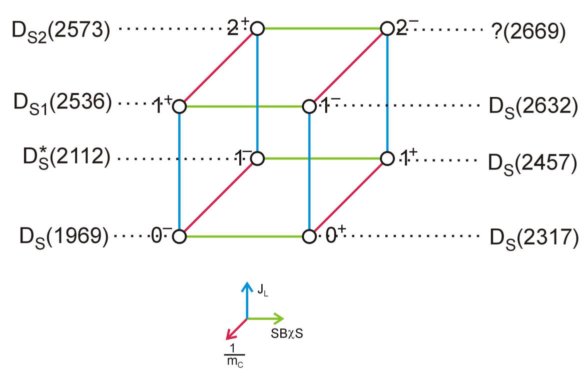

Fig. 4. Cube representing *schematic* (*e.g.* the units in the upper and lower plaquettes are different) classification of chiral doublers. Labels correspond to the case of $c\bar{s}$ mesons. Selex signal $D_s(2632)$ is interpreted as an excited doubler, see text.

The three-dimensional "cube" is aligned along three "directions":

— chiral symmetry breaking denoted by SBχB (horizontal)

— Isgur–Wise symmetry breaking $1/m_c$ (skew)

— total light angular momentum j_l (vertical).

We expect similar form of cartoon for $c\bar{u}$, $c\bar{d}$, $b\bar{s}$, $b\bar{u}$, $b\bar{d}$ mesons.

- First, we focus on $c\bar{s}$ mesons — $D_{\rm s}$ cube:

Lower left rung represents known pseudoscalar (0^-) $D_{\rm s}(1969)$ and vector (1^-) $D^\star_{\rm s}(2112)$, with $j_l = 1/2$ light angular momentum. The splitting between them (143 MeV) is an $1/m_c$ effect and is expected to vanish in infinitely heavy charm quark limit, *i.e.* both particles would have form the H multiplet. The upper left rung corresponds to $j_l = 3/2$ representation, *i.e.* 1^+ and 2^+ excited multiplet. Here $D_{\rm s1}(2536)$ and $D^\star_{{\rm s}J}(2573)$ are the candidates, separated by (smaller for excited states, here only 37 MeV) $1/m_c$ origin mass splitting. This "left plaquette" of the $D_{\rm s}$-cube represents well known states, before BaBar discoveries (we can call them "pre-BaBarian").

The novel aspect of the chiral doublers is the appearance of the *right plaquette*. First, we expect two chiral partners for $D_{\rm s}$ and $D^\star_{\rm s}$, representing right lower rung. Here newly discovered $D^\star_{{\rm s}J}(2317)$ and $D_{{\rm s}J}(2460)$ are the candidates for the $(0^+, 1^+)$ scalar-axial G multiplet. The averaged splitting for $(0^+, 0^-)$ and the averaged splitting for $(1^+, 1^-)$ are 349.2 ± 0.8 MeV and 346.8 ± 1.1 MeV, respectively, *i.e.* almost identical, as predicted a decade ago [15,16]. The splitting within the G multiplet, *i.e.* between the masses of the new BaBar state and CLEO state, is identical to the splitting between the $(1^-, 0^-)$ pair.

Let us move to upper light angular momentum (j_l direction in Fig. 2). We would also expect the chiral partners for the excited $j_l = 3/2$ multiplet, *i.e.* new chiral pair $(1^-, 2^-)$ [22]. Alternatively, this pair could be also viewed as the $j_l = 3/2$ excitation of the BaBar-CLEO $(0^+, 1^+)$ multiplet. The states within this new multiplet would be separated by similar $1/m_c$ split, like the split between $D_{\rm s2}$ and $D_{\rm s1}$, *i.e.* by 37 MeV. However, the question is how large is the chiral split for the excited states? Is it also equal to 350 MeV alike the chiral split for the $j_l = 1/2$ plaquette or is different? One can try to get some insight using the construction for effective chiral action for excited mesons [22]. Note that chiral shift for excited states is approximately half of the value of the shift for $j_l = 1/2$ multiplet (170 MeV).

The fact that excited states are less sensitive to the effects of the QCD vacuum is not totally unexpected, see *e.g.* [23]. Of course, the precise value of the chiral shift for the excited doubler can be provided only by an experiment. It is tempting to speculate that the very recent signal reported by Selex [7] is a (1^-) doubler of $D_{\rm s1}$, if the state is confirmed and its spin-parity is indeed (1^-). Then the chiral shift for excited strange charmed mesons would be of the order of 100 MeV only. If indeed this is the case, a natural expectation in the chiral doubler scenario is the presence of the chiral doubler for $D_{\rm s2}$ state as well, *i.e.* one would expect new, 2^- state within few MeV around 2669 MeV, possibly in $D_{\rm s}^\star \eta$ channel, to follow the pattern of the decay of other doublers.

- Now, we consider non-strange charmed mesons (D-cube):

Left plaquette is formed by known non-strange charmed mesons, *i.e.* for $j_\ell = 1/2$ we have pseudoscalar $D(1865)$ and vector $D^\star(2010)$, excited multiplet ($j_\ell = 3/2$) is formed by (1^+) $D_1(2420)$ and (2^+) $D_2(2460)$. Here two states from Belle, $D_0^\star(2308)$ and $D_1'(2427)$ are natural candidates for lower right rung of the D-cube, *i.e.* for the chiral doublers of $D(1865)$ and $D^\star(2010)$. There are however broad, since neither kinematic nor isospin restrictions apply here, contrary to their strange cousins. The precise value of the chiral shift is still an open problem, due to the experimental errors and systematic difference between the FOCUS [6] and Belle [5] signals. We would like to mention, that the fact that chiral mass shift seems to be equal of even larger for the non-strange mesons than for the strange ones, is not in contradiction with certain models of spontaneous breakdown of the chiral symmetry [17].

- Let us mention for completeness about bottom mesons ($B_{\rm s}$ and B):

In this case chiral doubling should be more pronounced, since the $1/m_Q$ corrections are three times smaller, *i.e.* the skew-symmetric (red) edges of the cubes are three times shorter, for $j_\ell = 1/2$ and $j_\ell = 3/2$ states, correspondingly. For $m_{\rm s} = 150$ MeV, we expect the chiral partners of $B_{\rm s}$ and $B_{\rm s}^\star$ to be 323 MeV heavier, while the chiral partners of B and $B^\star$ to be 345 MeV heavier [17]. We note that any observation of chiral doubling for B mesons would be a strong validation for chiral doublers proposal. For several recently proposed alternative scenarios for new states (multiquark states, hadronic molecules, modifications of quark potential, unitarization) a repeating pattern from charm to bottom seems to be hard to achieve without additional assumptions.

4. Baryons

In this lecture we discuss briefly the possibility of an extension of the chiral doublers scenario for all heavy–light baryons, including the exotic states like pentaquarks. To avoid any new parameters, we simply view baryons as solitons of the effective mesonic Lagrangian including both chiral copies of heavy–light mesons, a point addressed already in [15] and recently reanalyzed in [24]. We are working in large N_c limit, which justifies the soliton (Skyrmion) picture, and large heavy quark mass limit, where we exploit the Isgur–Wise symmetry. This approach could be viewed as a starting point for including $1/m_h$ corrections from the finite mass of the heavy quark, explicit breaking of chiral symmetry, *etc.*, alike the presented previously scheme does it for the mesons.

We follow here the scheme mentioned in the textbook [25], a variant of the original work by [26]. Charmed hyperons emerge as bound states of D and D^* in the presence of the SU(2) Skyrme background. First the pseudoscalar-vector heavy meson pair is being bound in the background of the static soliton, generating the $O(N_c^0)$ binding. Vibrational modes are the "fast degrees" of the freedom. The adiabatical rotation of the bound system by quantization of collective coordinates of the SU(2) Skyrmions alike proposed by Witten [27] corresponds then to "slow degrees" of freedom. It is well known, that in this case the rotation is not the free one. Fast degrees of freedom in Born–Oppenheimer approximation generate the effective "gauge" potential, of a Berry phase [28] type. In the case of degenerate pseudoscalar and vector mesons (IW limit) the phases coming from D meson and D^* meson are equal, but opposite. Their cancellation corresponds to the realization of the Isgur–Wise symmetry at the baryonic level, therefore degeneration of spin 1/2 and 3/2 multiplets. The details of this approach were outlined in [24]. The difference in respect to other similar works in the literature [9, 12, 26, 29] was to consider the full heavy–light effective Lagrangian with both chiral copies [15, 16] and to include the crucial effects of the chiral shift. We can imagine four different scenarios:

- Soliton of the light sector with baryon number 1 binds the H-multiplet — the resulting bound states exhibits the quantum numbers of the charmed baryons with standard $1/2^+$ parity.

- Soliton of the light sector with baryon number 1 binds the G-multiplet — the resulting bound states exhibits the quantum numbers of the charmed baryons with opposite $1/2^-$ parity.

- Soliton of the light sector with baryon number 1 binds the anti-flavored $\bar{H}$-multiplet — the resulting bound states exhibits the quantum numbers of the charmed baryon with minimal content of *five quarks* with standard $1/2^+$ parity, *i.e.* charmed pentaquark.

- Soliton of the light sector with baryon number 1 binds the anti-flavored $\bar{G}$-multiplet — the resulting bound states exhibits the quantum numbers of the charmed baryon with minimal content of *five quarks* with opposite $1/2^-$ parity, *i.e.* the chiral partner of the pentaquark.

Let us consider now the full mesonic effective Lagrangian defined in the previous section. First we observe, that due to the properties of the heavy spin symmetry, one can trade $\gamma^\mu A_\mu$ into $v^\mu A_\mu$ in the mixed term involving H, G and the axial current. This implies, that in the rest frame *static* Skyrmion background decouples the G and H Lagrangians. This decoupling allows immediately to write down the generic mass formula for opposite parity partner of the isoscalar baryon and for opposite parity partner of the isoscalar pentaquark (denoted by tilde)

$$\begin{aligned} \tilde{M} &= M_{\rm sol} + m_{\tilde{D}} - 3/2 g_G F'(0) + 3/(8I_1)\,, \\ \tilde{M}_5 &= M_{\rm sol} + m_{\tilde{D}} - 1/2 g_G F'(0) + 3/(8I_1) \end{aligned} \tag{32}$$

in analogy to identical formulae for the known sector for H, with D mesons and g_H axial couplings, respectively. The ordering of mass terms is as follows: first term corresponds to classical mass of the soliton (of order N_c), second term measures the (model-dependent via the shape of the soliton profile $F(r)$) binding with respect to the mass of the meson (independent on the number of colors) and the last term measures the $1/N_c$ split due to the moment of inertia I_1 of the soliton. It is of primary importance that both Hamiltonians for H and G sectors have the same functional form of lowest eigenvalue: M_5 for H and $\tilde{M}_5$ for G. Hence both parity partners emerge as H and G bound states in the SU(2) solitonic background. The mass difference comes in the first approximation solely from the difference of the coupling constants $g_G - g_H$ and meson mass difference $m_{\tilde{D}} - m_D$ where $m_{\tilde{D}} = (3M_{\tilde{D}*} + M_{\tilde{D}})/4$ is the averaged over heavy-spin mass of the $(1^+, 0^+)$ mesons. Constant g_G is the axial coupling constant in the opposite parity channel, responsible for pionic decays of the 1^+ axial states into 0^+ scalars. Using recent Belle data [3], *i.e.* 0^+ candidate D_0^* ($2308{\pm}17{\pm}15{\pm}28$) and 1^+ candidate D_1' ($2427 \pm 26 \pm 20 \pm 17$), we get $M_{\tilde{D}} = 2397$ MeV, unfortunately with still large errors.

One can easily combine the formulae for four, above mentioned, generic scenarios. Fist, we notice, that the mass splitting between the usual baryons of opposite parity leads to

$$\Delta_B = \Delta_M + 3/2 F'(0) g_H \delta g\,, \tag{33}$$

where $\Delta_M = M_{\tilde{D}} - M_D$ is the mass shift between the opposite parity heavy–light mesons and $\delta_g = 1 - g_G/g_H$ measures the difference between the axial couplings for both copies. Similar reasoning leads to the formula for the parity splitting between the opposite parity pentaquarks:

$$\Delta_P = \Delta_M + 1/2 F'(0) g_H \delta g\,. \tag{34}$$

Combining both formulae we get

$$\Delta_P = \frac{\Delta_B + 2\Delta_M}{3} \,. \tag{35}$$

Let us turn now towards the data. Comparing the mass shift between the lowest Λ_c states of opposite parities, $\Lambda_c(1/2^+,\ 2285)$ and $\Lambda_c(1/2^-,\ 2593)$ we arrive at $\Delta_B = 310$ MeV. Similarly, $\Xi_c(1/2^+,\ 2470)$ and $\Xi_c(1/2^-,\ 2790)$ give $\Delta_B = 320$ MeV. Comparing the shift of the opposite parity heavy charmed mesons from very recent Belle [5] data we arrive at $\Delta_M = 425$ MeV unfortunately with still large errors. These two numbers allow us to estimate $\Delta_P = 350$ MeV ± 60 MeV, *i.e.* we get the mass of the chiral doubler of the pentaquark as high as 3052 ± 60 MeV. We note that the argument proposed here is based on the leading approximation in large N_c and large m_h limit, and is intended to demonstrate the order of magnitude for parity splitting for heavy pentaquarks. Let us contrast these predictions to others in the literature (see the Table).

TABLE

Predicted masses of charmed pentaquark Θ_c^0 $(udud\bar{c})$

Model	Mass [MeV]	Ref.
constituent quark model (FS)	2902	[30]
diquark model	2710	[10]
diquark–triquark model	2985 ± 50	[11]
chiral soliton model	2704	[9, 12]
chiral doublers scenario	2700 ; 3052 ± 60	[24]
lattice calculation	2977	[31]

One is therefore tempted to interpret the recent H1 state [8] as a parity partner $\tilde{\Theta}_c$ of the yet undiscovered isosinglet pentaquark Θ_c of opposite parity and $M_5 \approx 2700$ MeV. Similar reasoning applies to other isospin channels, strange charmed pentaquarks and to extensions for b quarks. Despite BaBar and CLEO data yield with the impressive accuracy the chiral mesonic shift to be equal to 350 MeV, no charmed strange baryon data for both parities do exist by now, so one cannot make similar estimation for strange charmed pentaquarks.

5. Conclusions

In these lectures, we presented the basic concepts of the chiral doublers scenario. Our analysis, since designed as an introductory lectures, was to large extend qualitative and was based on the idealized leading approximation in the mass of the heavy quark mass. In the real world, "heavy" masses are finite, the effects of explicit breakdown of the light quark are present, chiral corrections should be organized in a systematic way, coupling to external (electroweak) currents should be present .... All this requires an extensive, systematic analysis, with many unknown till today parameters for the couplings in the subleading terms in the interaction. On the other side, such systematic analysis offers the possibility of honest, quantitative confrontation of the theoretical predictions with the experimental data. We would like to mention, that since the backbone of the chiral doublers scheme consists of the pattern of several symmetries of the QCD, the form of the possible interactions is quite constrained — *i.e.* the scheme is easy to falsify or verify, provided sufficient amount of experimental data will be available. It is encouraging, that very recent realistic study of electromagnetic decays of chiral properties based on subleading terms [32] is in general consistent with chiral doublers scenario and seems to favor this interpretation of new narrow charmed mesons in comparison to "hadronic molecules" scenario. We do hope that these lectures will encourage further serious investigations of the possibility of chiral doublers scenario both for heavy–light mesons and for heavy–light baryons.

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REFERENCES

[1] BABAR Coll., B. Aubert *et al.*, *Phys. Rev. Lett.* **90**, 242001 (2003) [arXiv:hep-ex/0304021].

[2] CLEO Coll., D. Besson *et al.*, *Phys. Rev.* **D68**, 032002 (2003) [arXiv:hep-ex/0305100].

[3] Belle Coll., P. Krokovny *et al.*, *Phys. Rev. Lett.* **91**, 262002 (2003) [arXiv:hep-ex/0308019].

[4] BABAR Coll., B. Aubert *et al.*, *Phys. Rev.* **D69**, 031101 (2004) [arXiv:hep-ex/0310050].

[5] Belle Coll., K. Abe *et al.*, *Phys. Rev.* **D69**, 112002 (2004) [hep-ex/0307021].

[6] FOCUS Coll., E.W. Vaandering *et al.*, hep-ex/0406044.

[7] SELEX Coll., A.V. Evdokimov *et al.*, `hep-ex/0406045`.

[8] H1 Coll., `hep-ex/0403017`.

[9] Y. Oh, B.-Y. Park, D.-P. Min, *Phys. Lett.* **B331**, 362 (1994) [`arXiv:hep-ph/9405297`]; *Phys. Rev.* **D50**, 3350 (1994) [`arXiv:hep-ph/9407214`].

[10] R.L. Jaffe, F. Wilczek, *Phys. Rev. Lett.* **91**, 232003 (2003) [`arXiv:hep-ph/0307341`].

[11] M. Karliner, H.J. Lipkin, `arXiv:hep-ph/0307343`.

[12] B. Wu, B.-Q. Ma, *Phys. Rev.* **D70**, 034025 (2004) [`hep-ph/0402244`].

[13] see e.g.: R.N. Cahn, J.D. Jackson, *Phys. Rev.* **D68**, 037502 (2003); T. Barnes, F.E. Close, H.J. Lipkin, *Phys. Rev.* **D68**, 054006 (2003); H.-Y. Cheng, W.-S. Hou, *Phys. Lett.* **B566**, 193 (2003); A.P. Szczepaniak, *Phys. Lett.* **B567**, 23 (2003); S. Godfrey, *Phys. Lett.* **B568**, 254 (2003); G.S. Bali, *Phys. Rev.* **D68**, 071501 (2003); S. Nussinov, `hep-ph/0306187`; Y.-B. Dai, C.-S. Huang, C. Liu, S.-L. Zhu, *Phys. Rev.* **D68**, 114011 (2003); K. Terasaki, *Phys. Rev.* **D68**, 011501 (2003) [`hep-ph/0305213`]; A. Dougall, R.D. Kenway, C.M. Maynard, C. McNelie, *Phys. Lett.* **B569**, 41 (2003); T.E. Browder, S. Pakvasa, A.A. Petrov, *Phys. Lett.* **B578**, 365 (2004); A. Deandrea, G. Nardulli, A.D. Polosa, *Phys. Rev.* **D68**, 097501 (2003); Ch.-H. Chen, H.N. Li, *Phys. Rev.* **D69**, 054002 (2004); M. Sadzikowski, *Phys. Lett.* **B579**, 39 (2004); A. Datta, P.J. O'Donnell, *Phys. Lett.* **B572**, 164 (2003); M. Suzuki, `hep-ph/0307118`; P. Bicudo, `hep-ph/0401106`; E. van Beveren, G. Rupp, *Phys. Rev. Lett.* **91**, 012003 (2003); *Eur. Phys. J.* **C32**, 493 (2004); M.F.M. Lutz, E.E. Kolomeitsev, *Nucl. Phys.* **A730**, 392 (2004); E.E. Kolomeitsev, M.F.M. Lutz, *Phys. Lett.* **B582**, 39 (2004); M. Harada, M. Rho, Ch. Sasaki, *Phys. Rev.* **D70**, 074002 (2004) [`hep-ph/0312182`]; I. Maiani, F. Piccinini, A.D. Polosa, V. Riquer, *Phys. Rev.* **D70**, 054009 (2004) [`hep-ph/0407025`]; B. Nicolescu, J.P.B.C. de Melo, `hep-ph/0407088`; Yu-Qi Chen, Xue-Qian Li, `hep-ph/0407062`; R. Ferrandes, `hep-ph/0407212`.

[14] For recent review, see F. Colangelo, F. De Fazio, R. Ferrandes, *Mod. Phys. Lett.* **A19**, 2083 (2004) [`hep-ph/0407137`].

[15] M.A. Nowak, M. Rho, I. Zahed, *Phys. Rev.* **D48** 4370 (1993).

[16] W.A. Bardeen, C.T. Hill, *Phys. Rev.* **D49**, 409 (1994).

[17] W.A. Bardeen, E.J. Eichten, C.T. Hill, *Phys. Rev.* **D68**, 054024 (2003); M.A. Nowak, M. Rho, I. Zahed, *Acta Phys. Pol. B* **35**, 2377 (2004) [`hep-ph/0307102`]; M. Harada, M. Rho, C. Sasaki, *Phys. Rev.* **D70**, 074002 (2004) [`hep-ph/0312182`]; P. Colangelo, F. De Fazio, *Phys. Lett.* **B570**, 180 (2003).

[18] N. Isgur, M.B. Wise, *Phys. Rev. Lett.* **66**, 1130 (1991). For early suggestions, see also E.V. Shuryak, *Phys. Lett.* **B93**, 134 (1980); *Nucl. Phys.* **B198**, 83 (1982).

[19] H. Georgi, *Phys. Lett.* **B240**, 447 (1990).

[20] See, *e.g.*, I. Zahed, G. Brown, *Phys. Rep.* **142**, 1 (1986); R.D. Ball, "Desperately seeking mesons", in *Proc. of the Workshop on Skyrmions and Anomalies*, ed. by M. Jeżabek and M. Praszałowicz, p. 54, World Scientific 1987.

[21] M.B. Wise, *Phys. Rev.* **D45**, R2118 (1992); T.-M. Yan *et al.*, *Phys. Rev.* **D46**, 1148 (1992); G. Burdman, J.F. Donoghue, *Phys. Lett.* **B280**, 287 (1992).

[22] M.A. Nowak, I. Zahed, *Phys. Rev.* **D48**, 356 (1993).

[23] L.Ya. Glozman, `hep-ph/0309334`.

[24] M.A. Nowak, M. Praszałowicz, M. Sadzikowski, J. Wasiluk, *Phys. Rev.* **D70**, 031503 (2004) [`arXiv:hep-ph/0403184`].

[25] M.A. Nowak, M. Rho, I. Zahed, *Chiral Nuclear Dynamics*, World Scientific, Singapore 1995, pp. 355, 362.

[26] E. Jenkins, A.V. Manohar, M.B. Wise, *Nucl. Phys.* **B396**, 27 (1993) [`arXiv:hep-ph/9205243`]; Z. Guralnik, M.E. Luke, A.V. Manohar, *Nucl. Phys.* **B390**, 474 (1993) [`arXiv:hep-ph/9208221`].

[27] E. Witten, *Nucl. Phys.* **B223** 422, 433 (1983).

[28] F. Wilczek, A. Zee, *Phys. Rev. Lett.* **52**, 2111 (1984).

[29] M. Harada, F. Sannino, J. Schechter, H. Weigel, *Phys. Rev.* **D56**, 4098 (1997) [`arXiv:hep-ph/9704358`].

[30] Fl. Stancu, `arXiv:hep-ph/0410033`.

[31] T.-W. Chiu, T.-H. Hsieh, `arXiv:hep-ph/0404007`.

[32] T. Mehen, R.P. Springer, *Phys. Rev.* **D70**, 074014 (2004) [`hep-ph/0407181`].