SU(6) extension of the Weinberg-Tomozawa meson-baryon Lagrangian

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A consistent SU(6) extension of the Weinberg-Tomozawa meson-baryon chiral Lagrangian is constructed which incorporates vector meson and baryon decuplet degrees of freedom. The corresponding Bethe-Salpeter approximation predicts the existence of an isoscalar spin-parity $\frac{3}{2}^-$ K^*N bound state (strangeness $+1$) with a mass around 1.7–1.8 GeV. It is the highest hypercharge state of an antidecuplet SU(3) representation and it is unstable through K^* decay. The estimated width of this state (neglecting d-wave KN decay) turns out to be small ($\Gamma \leq 15$ MeV). Clear signals of this resonance would be found in reactions like $\gamma p \rightarrow \bar{K}^0 pK^+\pi^-$ by looking at the three body $pK^+\pi^-$ invariant mass.

I. INTRODUCTION

Forty years ago, it was suggested [1–3] that it might be a useful approximation to assume that the light quark-light quark interaction is approximately spin independent as well as SU(3) independent. This corresponds to treating the six states of a light quark (u, d or s with spin up, ↑, or down, ↓) as equivalent, and leads us to the invariance group SU(6). In order that we can speak meaningfully of SU(6) transformations affecting spin but not orbital angular momentum (L) as invariances, it must be assumed that the orbital angular momentum and the quark spin are to a good approximation, separately conserved. This, in turn requires the spin-orbit, tensor and spin-spin interactions between quarks to be small. As is known, mixing the compact, purely internal flavor symmetry, with the noncompact Poincare symmetry of spin angular momentum led to some inconsistencies, which gave rise to the no-go Coleman-Mandula theorem [4] forbidding such exact hybrid symmetries, unless supersymmetry is invoked. However, there exist several SU(6) predictions (relative closeness of baryon octet and decuplet masses, the axial current coefficient ratio $F/D = 2/3$, the magnetic moment ratio $\mu_p/\mu_n = -3/2$) which are remarkably well satisfied in nature [5]. This suggests that SU(6) could be a good approximate symmetry. Indeed, in the large N_c limit (being N_c the number of colors) [6,7], there exists an exact spin-flavor symmetry for ground state baryons [8]. Moreover, though in general the spin-flavor symmetry is not exact for excited baryons even in the large N_c limit, in the real world ($N_c = 3$), the zeroth order spin-flavor symmetry breaking turns out to be similar in magnitude to $\mathcal{O}(N_c^{-1})$ breaking effects [9]. In the meson sector, an underlying static chiral $U(6) \times U(6)$ symmetry has been advocated by Caldi and Pagels [10,11], in which vector mesons would be "dormant" Goldstone bosons acquiring mass through relativistic corrections. This scheme solves a number of theoretical problems in the classification of mesons and also makes predictions which are in remarkable agreement with the experiment. In any case, although spin-flavor symmetry in the meson sector is not a direct consequence of large N_c QCD, vector mesons (K^*, ρ, ω, $\bar{K}^*$, ϕ) do exist and they are known to play a relevant role in hadronic physics [12]. Inescapably, they will couple to baryons and will presumably influence the properties of the baryonic resonances. Lacking better theoretically founded models to include vector mesons, we regard the spin-flavor symmetric scenario as a reasonable first step. The large N_c consequences of this scheme have been pursued in [13].

Since the pure SU(3) (flavor) transformations commute with the pure SU(2) (spin) transformations, it follows that a SU(6) multiplet can be decomposed into SU(3) multiplets each of definite total spin. With the inclusion of spin there are 36 quark-antiquark ($q\bar{q}$) states, and the SU(6) group representation reduction (denoting the SU(6) multiplets by their dimensionality and a SU(3) multiplet μ of spin J by μ_{2J+1}) reads

$$\mathbf{6} \otimes \mathbf{6}^* = \mathbf{35} \oplus \mathbf{1} = (\mathbf{8}_1 \oplus \mathbf{8}_3 \oplus \mathbf{1}_3) \oplus \mathbf{1}_1. \quad (1)$$

We might expect the lowest bound state to be a s-state and since the relative parity of a fermion-antifermion pair is odd, the octet of pseudoscalar (K, π, η, $\bar{K}$) and the nonet of vector (K^*, ρ, ω, $\bar{K}^*$, ϕ) mesons are commonly placed in the **35** representation. Note that the **35** allows nine vector mesons but only eight 0^- mesons. A ninth 0^- meson (η') must go in the **1** of SU(6). This may account for the phenomenological evidence that the mixing of the octet and singlet states is much smaller for the 0^- mesons than for the 1^- mesons [14]. Mesons of spin greater than one can be understood as states of the $q\bar{q}$ system with $L > 0$. On the other hand, with the inclusion of the spin there are 216 three quark states, and it follows

$$\mathbf{6} \otimes \mathbf{6} \otimes \mathbf{6} = \mathbf{56} \oplus \mathbf{20} \oplus \mathbf{70} \oplus \mathbf{70}$$

$$= (\mathbf{8}_2 \oplus \mathbf{10}_4) \oplus (\mathbf{1}_4 \oplus \mathbf{8}_2)$$

$$\oplus 2 \times (\mathbf{10}_2 \oplus \mathbf{8}_4 \oplus \mathbf{8}_2 \oplus \mathbf{1}_2). \quad (2)$$

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It is natural to assign the lowest-lying baryons to the **56** of SU(6), since it can accommodate an octet of spin-1/2 baryons and a decuplet of spin-3/2 baryons, which are precisely the SU(3)-spin combinations of the low-lying baryon states $((N, \Sigma, \Lambda, \Xi)$ and $(\Delta, \Sigma^*, \Xi^*, \Omega))$. Furthermore, the **56** of SU(6) is totally symmetric, which allows the baryon to be made of three quarks in *s*-wave (the color wavefunction being antisymmetric).

Here we will consider describing the *s*-wave interaction between the lowest-lying meson multiplet (**35**) and the lowest-lying baryons (**56**-plet) at low energies. At larger energies higher partial waves are involved and a suitable treatment of spin-orbit effects in the SU(6) scheme would be required. Thus, assuming that the *s*-wave effective meson-baryon Hamiltonian is SU(6) invariant, and since the SU(6) decomposition of the product of the **35** (meson) and **56** (baryon) representations yields

$$35 \otimes 56 = 56 \oplus 70 \oplus 700 \oplus 1134, \quad (3)$$

we have only four (Wigner-Eckart irreducible matrix elements) free functions of the meson-baryon Mandelstam variable *s*. Similar ideas were already explored in the late sixties, within the effective range approximation [15]. Here, we introduce two major improvements: (i) We make use of the underlying Chiral Symmetry (CS), which allows us to determine the value of the SU(6) irreducible matrix elements from the Weinberg-Tomozawa (WT) interaction [16,17], the leading term of the chiral Lagrangian involving Goldstone bosons and the octet of spin-1/2 baryons. This is not a trivial fact and it is intimately linked to the underlying group structure of the WT term. (ii) We go beyond the effective range approximation and follow a scheme based on the solution of the Bethe-Salpeter-Equation (BSE), which incorporates two-body coupled channel unitarity and has been successfully employed in the study of *s*-wave $(\mathbf{8}_1)$meson-$(\mathbf{8}_1)$meson and $(\mathbf{8}_1)$meson-$(\mathbf{8}_2, \mathbf{10}_4)$baryon scattering and resonances, within different renormalization schemes [18–28].

II. SU(6) MESON-BARYON EFFECTIVE INTERACTION MATRIX

We will work with well defined total isospin (I), angular momentum (J) and hypercharge (Y) (strangeness plus baryon numbers) meson-baryon states constructed out of the SU(6) **35** (mesons) and **56** (baryon) multiplets. In what follows we always use the labels μ and ϕ to denote generic SU(3) and SU(6) representations, respectively. For short, we use the notation $\mathcal{M} \equiv [(\mu_M)_{2J_M+1}, I_M, Y_M]$ for mesons and similarly for baryons ($\mathcal{B}$). Thus, $\mu_M = \mathbf{8}, \mathbf{1}$ and $\mu_B = \mathbf{8}, \mathbf{10}$ are the meson and baryon SU(3) multiplets, respectively, and $J_{M,B}, I_{M,B}, Y_{M,B}$ are the spin, isospin and hypercharge quantum numbers of the involved hadrons. The meson-baryon states in terms of the SU(6) coupled (orthonormal) basis read

$$|\mathcal{MB}; JIY\rangle = \sum_{\mu,\alpha,\phi} \begin{pmatrix} \mu_M & \mu_B & \mu \\ I_M Y_M & I_B Y_B & IY \end{pmatrix}$$
$$\times \begin{pmatrix} \mathbf{35} & \mathbf{56} & \phi \\ \mu_M J_M & \mu_B J_B & \mu J \alpha \end{pmatrix} |\phi; \mu^\alpha_{2J+1} IY\rangle, \quad (4)$$

where $Y = Y_M + Y_B$, $|I_M - I_B| \leq I \leq I_M + I_B$, and for *s*-wave scattering $|J_M - J_B| \leq J \leq J_M + J_B$, while $\phi = $ **56**, **70**, **700**, **1134**, and α accounts for the multiplicity of each of the μ_{2J+1} SU(3) multiplets of spin J (for $L = 0$, J is given by the total spin of the meson-baryon system) entering in the representation ϕ. Multiplicities higher than one only happen for the **1134** representation, where the $\mathbf{27}_4$, $\mathbf{10}_4$, $\mathbf{27}_2$ and the $\mathbf{10}_2$ multiplets appear twice and the $\mathbf{8}_4$ and $\mathbf{8}_2$ ones appear 3 times. The index μ runs over the $(\mathbf{27}, \mathbf{10}, \mathbf{10}^*, \mathbf{8}_s, \mathbf{8}_a, \mathbf{1})$, $(\mathbf{35}, \mathbf{27}, \mathbf{10}, \mathbf{8})$, $(\mathbf{8})$ and $(\mathbf{10})$ SU(3) representations for the octet-octet, octet-decuplet, singlet-octet and singlet-decuplet decompositions, respectively. Finally in Eq. (4), the two coefficients which multiply each element of the SU(6) coupled basis are the SU(3) isoscalar factors [29], and the SU(6)-multiplet coupling factors [30,31]. The assumption that the effective *s*-wave meson-baryon *potential* (V) is a SU(6) invariant operator implies that (i) the coupled states $|\phi; \mu^\alpha_{2J+1} IY\rangle$ are eigenvectors of V and (ii) the corresponding eigenvalues $V_\phi(s)$ may depend on the SU(6) representation ϕ but not on the other quantum numbers μ, α, J, I, or Y. Thus, in the noncoupled basis we find

$$\langle \mathcal{M'B'}; JIY|V|\mathcal{MB}; JIY\rangle = \sum_\phi V_\phi(s) \mathcal{P}^{\phi, JIY}_{\mathcal{MB}, \mathcal{M'B'}}, \quad (5)$$

with projectors given by

$$\mathcal{P}^{\phi, JIY}_{\mathcal{MB}, \mathcal{M'B'}} = \sum_{\mu,\alpha} \begin{pmatrix} \mathbf{35} & \mathbf{56} & \phi \\ \mu_M J_M & \mu_B J_B & \mu J \alpha \end{pmatrix} \begin{pmatrix} \mu_M & \mu_B & \mu \\ I_M Y_M & I_B Y_B & IY \end{pmatrix}$$
$$\times \begin{pmatrix} \mu'_{M'} & \mu'_{B'} & \mu \\ I'_{M'} Y'_{M'} & I'_{B'} Y'_{B'} & IY \end{pmatrix}$$
$$\times \begin{pmatrix} \mathbf{35} & \mathbf{56} & \phi \\ \mu'_{M'} J'_{M'} & \mu'_{B'} J'_{B'} & \mu J \alpha \end{pmatrix}. \quad (6)$$

III. CHIRAL SYMMETRY CONSTRAINTS

We make use of the underlying CS and propose a chiral expansion to determine the $V_\phi(s)$ functions. Thus, we look at the effective *s*-wave *potential* describing the interaction of the Goldstone pseudoscalar meson and the lowest $J^P = \frac{1}{2}^+$ baryon octets. From the SU(3) WT chiral Lagrangian (we use the convention $V = -\mathcal{L}$), one finds for each (I, Y) sector $[(0, 2), (1, 2), (1/2, 1), (3/2, 1), (0, 0), (1, 0), (2, 0), (1/2, -1), (3/2, -1), (0, -2), (1, -2)]$ and on the mass shell (recall that $J = 1/2$)

$$V^{IY}_{ab}(\sqrt{s}) = D^{IY}_{ab} \frac{2\sqrt{s} - M_a - M_b}{4f^2} \quad (7)$$

with

$$D^{IY} = \sum_{\mu,\gamma,\gamma'} \lambda_{\mu_\gamma \to \mu_{\gamma'}} \begin{pmatrix} \mathbf{8} & \mathbf{8} \\ I_M Y_M & I_B Y_B \end{pmatrix} \begin{vmatrix} \mu_\gamma \\ IY \end{vmatrix}$$
$$\times \begin{pmatrix} \mathbf{8} & \mathbf{8} \\ I'_{M'} Y'_{M'} & I'_{B'} Y'_{B'} \end{pmatrix} \begin{vmatrix} \mu_{\gamma'} \\ IY \end{vmatrix}, \tag{8}$$

where M_b (M_a) is the baryon mass of the initial (final) channel, $f \simeq 93$ MeV the pion weak decay constant, μ runs over the **27**, **10**, **10***, **8** and **1** SU(3) representations and γ, γ' are used to account for the two octets ($\mathbf{8}_s$ and $\mathbf{8}_a$) which appear in the $\mathbf{8} \otimes \mathbf{8}$ decomposition ($\mathbf{27} + \mathbf{10} + \mathbf{10^*} + \mathbf{8}_s + \mathbf{8}_a + \mathbf{1}$). Besides, $\lambda_{\mathbf{27}} = 2$, $\lambda_{\mathbf{8}_s} = \lambda_{\mathbf{8}_a} = -3$, $\lambda_{\mathbf{1}} = -6$, $\lambda_{\mathbf{10}} = \lambda_{\mathbf{10^*}} = \lambda_{\mathbf{8}_s \leftrightarrow \mathbf{8}_a} = 0$, which reproduces the D-matrix eigenvalues found in Ref. [26]. Thus, we see that CS at leading order is much more predictive than SU(3) symmetry, and it predicts the values of the seven λ couplings, which otherwise will be totally arbitrary functions of s. Note that the SU(3) WT Lagrangian also provides the s dependence (($\sqrt{s} - M$), with M the common mass of the baryon octet in the SU(3) limit) of the effective potential, and thus one is left with only two free parameters, namely, f and M.

It is clear that not all SU(3) invariant interactions in the ($\mathbf{8}_1$)meson-($\mathbf{8}_2$)baryon sector can be extended to a SU(6) invariant interaction. Remarkably, the seven couplings (λ's) in the WT interaction turn out to be consistent with SU(6) and moreover, the extension is unique. In other words, there is a choice of the *four* couplings for the $\mathbf{35} \otimes \mathbf{56}$ interaction that, when restricted to the $\mathbf{8}_1 \otimes \mathbf{8}_2$ sector, reproduces the *seven* SU(3) WT couplings and such choice is unique. Indeed, the *potential* of Eq. (7) can be recovered, in the SU(3) limit, from Eq. (5) by taking

$$V_\phi(s) = \bar{\lambda}_\phi \frac{\sqrt{s} - M}{2f^2}, \tag{9}$$

with $\bar{\lambda}_{\mathbf{56}} = -12$, $\bar{\lambda}_{\mathbf{70}} = -18$, $\bar{\lambda}_{\mathbf{700}} = 6$ and $\bar{\lambda}_{\mathbf{1134}} = -2$ and M now being the common octet and decuplet baryon mass.[1] The underlying reason for this is CS. Indeed, the WT Lagrangian is not just SU(3) symmetric but also chiral ($SU_L(3) \otimes SU_R(3)$) invariant. Symbolically (and up to an overall coefficient)

$$\mathcal{L}_{WT} = \text{Tr}([M^\dagger, M]B^\dagger B). \tag{10}$$

This structure, dictated by CS, is more suitably analyzed in the t-channel. The mesons M fall in the representation **8**

which is also the adjoint representation. The commutator $[M^\dagger, M]$ indicates a t-channel coupling to the $\mathbf{8}_a$ (antisymmetric) representation, thus

$$\mathcal{L}_{WT} = ((M^\dagger \otimes M)_{\mathbf{8}_a} \otimes (B^\dagger \otimes B)_{\mathbf{8}})_{\mathbf{1}}. \tag{11}$$

The unique SU(6) extension is then

$$\mathcal{L}_{WT,SU(6)} = ((M^\dagger \otimes M)_{\mathbf{35}_a} \otimes (B^\dagger \otimes B)_{\mathbf{35}})_{\mathbf{1}}, \tag{12}$$

since the **35** is the adjoint representation of SU(6). The t-channel decompositions $\mathbf{35} \otimes \mathbf{35} = \mathbf{1} \oplus \mathbf{35}_s \oplus \mathbf{35}_a \oplus \mathbf{189} \oplus \mathbf{280} \oplus \mathbf{280^*} \oplus \mathbf{405}$ and $\mathbf{56} \otimes \mathbf{56^*} = \mathbf{1} \oplus \mathbf{35} \oplus \mathbf{405} \oplus \mathbf{2695}$ indicate that the coupling in Eq. (12) exists and is indeed unique, all coupling constants being reduced to a single independent one, namely, that of the WT Lagrangian (pion weak decay constant, besides the hadron masses).

The large N_c behavior of the SU(6) WT extension proposed here has been contemplated in [13]. Two interesting conclusions of that study are, (i) a consistent treatment of the WT interaction, yields a generic large N_c amplitude of the same order as the baryon pole one, $\mathcal{O}(N_c^0)$, in amplitudes involving vector mesons, instead of the known $1/N_c$ suppression for pseudoscalars. And (ii) the SU(6) WT interaction is large in the **70**-plet SU(6) sector even for large N_c, which is consistent with quark model considerations [32] (see also footnote [1]).

IV. MESON-BARYON SCATTERING MATRIX AND THE Θ^{*+} BARYON.

We solve the coupled channel BSE with an interaction kernel determined by Eqs. (5) and (9). Some mass breaking effects can be taken into account just by replacing ($\sqrt{s} - M$) by $(2\sqrt{s} - M_a - M_b)/2$ in Eq. (9). In a given JIY sector, the solution for the coupled channel s-wave scattering amplitude, $T_{IY}^J(\sqrt{s})$ (normalized as the t matrix defined in Eq. (33) of [24]), in the *on-shell* scheme [19–25] reads,

$$T_{IY}^J(\sqrt{s}) = \frac{1}{1 - V_{IY}^J(\sqrt{s})J_{IY}^J(\sqrt{s})} V_{IY}^J(\sqrt{s}) \tag{13}$$

with

$$V_{IY}^J(\sqrt{s}) = \langle \mathcal{M}'\mathcal{B}'; JIY|V|\mathcal{M}\mathcal{B}; JIY \rangle, \tag{14}$$

and $J_{IY}^J(\sqrt{s})$ a diagonal matrix of loop functions [24,25], which logarithmically diverge and hence need one subtraction or an ultraviolet (UV) cutoff to make them finite. Possible d-wave mixings (chiral corrections provide d-wave meson-baryon amplitudes) and the rich new phenomenology (there are new open channels, due to the inclusion of the vector meson degrees of freedom and the products from their decays, not taken into account in the usual meson-baryon analysis based on the WT interaction) which can be extracted from Eq. (13) will be studied elsewhere. We will focus here on the $Y = 2$ (strangeness $= +1$) sector of great interest since the claim by the LEPS

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[1]The SU(6) extension thus obtained (Eqs. (5) and (9)) also leads to the *potentials* used in Ref. [27,28] to study the ($\mathbf{10}_4$-)baryon-($\mathbf{8}_1$)meson sector. Note also that the **70** of SU(6) leads to the most attractive s-wave meson-baryon interaction. This would be consistent with the scenario studied in [9], namely, that the first negative parity baryon excited states are members of the **70** multiplet, and also with constituent quark model considerations, where the **70**-plet appears when one of the quarks occupies an orbitally excited level [32].

collaboration of the observation of the $\Theta^+(1540)$ resonance [33]. Though its existence is controversial (see e.g. [34]), it seems clear that the possible candidate would be an isoscalar, extremely narrow (with width definitely smaller than 10 MeV), while its spin-parity has not been established yet.

In the $Y = 2$, $I = 0$ sector, the $|KN\rangle$ and $|K^*N\rangle$ states appear. We have $|KN\rangle_{J=1/2} = (|\mathbf{700}; \mathbf{10}_2^*\rangle - \sqrt{3}|\mathbf{1134}; \mathbf{10}_2^*\rangle)/2$ and $|K^*N\rangle_{J=1/2} = (-\sqrt{3}|\mathbf{700}; \mathbf{10}_2^*\rangle - |\mathbf{1134}; \mathbf{10}_2^*\rangle)/2$, and $|K^*N\rangle_{J=3/2} = -|\mathbf{1134}; \mathbf{10}_4^*\rangle$. In the $J = 1/2$ channel, we find a resonance (pole in the second Riemann sheet [24,25]), though its exact position depends on the details of the Renormalization Scheme (RS) used. It comes out too wide ($\Gamma > 100$ MeV) to be identified as the $\Theta^+(1540)$. The situation is much more suggestive for $J = 3/2$. There, the effective interaction is determined by the **1134** SU(6) representation alone and it is attractive ($\bar{\lambda}_{\mathbf{1134}} = -2$). A pole is found in the first Riemann sheet corresponding to a K^*N bound state which we call Θ^{*+}. This state is unstable since the K^* decays into $K\pi$. In order to estimate the Θ^{*+} width, we model the Θ^*NK^* coupling as

$$\mathcal{L}_{\Theta^*NK^*} = -\frac{g}{\sqrt{2}} \bar{\Theta}^\mu (K_\mu^{*0} p - K_\mu^{*+} n) + \text{H.c.}, \quad (15)$$

while the K^* decay is described by

$$\begin{aligned}
\mathcal{L}_{K^*K\pi} = &-\frac{ig'}{\sqrt{2}f^2}\{\partial^\mu \pi^+ \partial^\nu K^-(\partial_\mu K_\nu^{*0} - \partial_\nu K_\mu^{*0}) \\
&+ \partial^\mu \pi^- \partial^\nu \bar{K}^0(\partial_\mu K_\nu^{*+} - \partial_\nu K_\mu^{*+}) \\
&- \frac{1}{\sqrt{2}}\partial^\mu \pi^0 \partial^\nu \bar{K}^0(\partial_\mu K_\nu^{*0} - \partial_\nu K_\mu^{*0}) \\
&+ \frac{1}{\sqrt{2}}\partial^\mu \pi^0 \partial^\nu K^-(\partial_\mu K_\nu^{*+} - \partial_\nu K_\mu^{*+})\} + \text{H.c.}
\end{aligned}$$
$$(16)$$

In the previous formulas, Θ^μ is a Rarita-Schwinger field, p and n the nucleon fields, K_μ^{*0}, $\bar{K}_\mu^{*0} = (K_\mu^{*0})^\dagger$ and $K_\mu^{*+} = (K_\mu^{*-})^\dagger$ the Proca fields which annhilate and create neutral and charged K^* and $\bar{K}^*$ mesons, and similarly for the kaon and antikaon fields, while π^0 and $\pi^+ = (\pi^-)^\dagger$ are the pion fields. The coupling g is determined by the residue at the pole of $T_{02}^{3/2}$ [i.e., $T_{02}^{3/2} \approx g^2 \times 2M_{\Theta^*}/(s - M_{\Theta^*}^2)$] and we fix $g' \approx 0.14$ to reproduce $\Gamma_{K^{*0}} \approx \Gamma_{K^{*+}} \approx 50$ MeV (we use charged averaged masses). This gives

$$\begin{aligned}
\Gamma_{\Theta^*} = &\frac{g^2 g'^2}{32\pi^3 f^4} \frac{1}{6M_{\Theta^*}} \int_{m_\pi+m_K}^{M_{\Theta^*}-M_N} d\tilde{m} \tilde{m}^2 \tilde{q} q_\pi^3 \\
&\times \frac{(2\tilde{m}^2 + \tilde{E}^2)(M_{\Theta^*} + M_N - \tilde{E})}{(\tilde{m}^2 - m_{K^*}^2)^2 + (\tilde{m}\tilde{\Gamma})^2},
\end{aligned} \quad (17)$$

where $\tilde{m}$, $\tilde{E}$, $\tilde{q}$, and $\tilde{\Gamma}$ are the invariant mass, energy, momentum, and width, respectively, of the virtual K^* in the at rest Θ^* system,
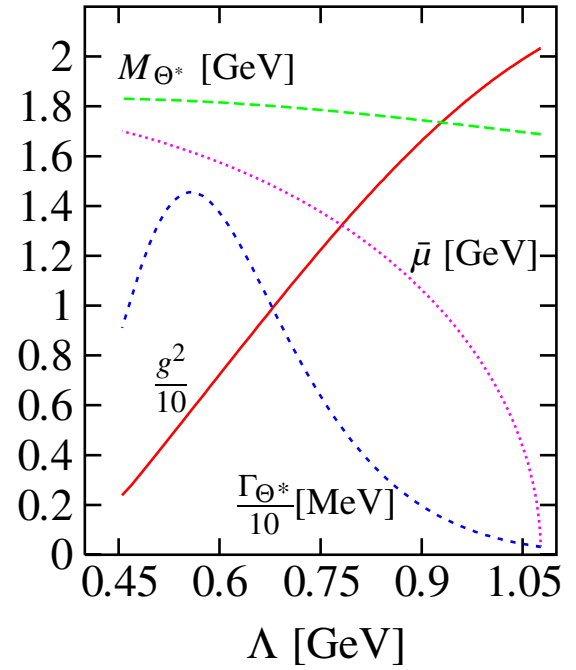

FIG. 1 (color online). Resonance Θ^{*+} properties as a function of the UV cutoff Λ or the subtraction scale $\bar{\mu}$.

$$\tilde{E} = \frac{M_{\Theta^*}^2 + \tilde{m}^2 - M_N^2}{2M_{\Theta^*}}, \qquad \tilde{q}^2 = \tilde{E}^2 - \tilde{m}^2,$$
$$\tilde{\Gamma} = \frac{g'^2}{32\pi f^4} q_\pi^3 \tilde{m}^2, \quad (18)$$

and q_π is the pion momentum in the at rest K^* system

$$q_\pi^2 = \frac{\lambda(\tilde{m}^2, m_\pi^2, m_K^2)}{(2\tilde{m})^2}. \quad (19)$$

Resonance mass, residue and width depend on the RS employed. We have used an UV cutoff (Λ) to evaluate the loop function $J(\sqrt{s})$, which is equivalent to choose an scale $\bar{\mu}$ such that $J(\sqrt{s} = \bar{\mu}) = 0$. Results are shown in Fig. 1. For $\bar{\mu}$ ranging from 0.05 GeV ($\Lambda \approx 1.08$ GeV) to 1.7 GeV ($\Lambda \approx 0.46$ GeV) the resonance mass (width) varies from 1.688 GeV (0.3 MeV), close to the ($M_N + m_\pi + m_K$) threshold, to 1.831 GeV (9 MeV, but the width does not grow monotonously, see figure), close to the ($M_N + m_{K^*}$) threshold.

In our treatment we have not included d-wave KN, K^*N contributions, nor further s-wave terms such as the u-channel pole graph or single pion exchange between K^* and N (the vertex $K^*K^*\pi$ being of abnormal parity) since none of these mechanisms contributes to the K^*N s-wave scattering length (i.e., they vanish at threshold).[2] Another mechanism for K^*N scattering not included here would be the sequential exchange of two pions with an

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[2]Except for a small contribution from the antibaryon term of the pole graph, all elementary vertices involved are p-wave.

intermediate K meson, corresponding to a box graph $K^*N \to KN \to K^*N$, which involves two p-wave normal parity $K^*K\pi$ vertices. Unlike the single pion exchange, such a contribution is not vanishing at threshold since the two virtual pions need not carry a small momentum. An analogous mechanism has been considered long ago for KN scattering [35,36], this time with the box graph $KN \to K^*N \to KN$. Technically a box graph is difficult to work with, since one must somehow renormalize its intrinsic UV divergence, and then renormalize its contribution in the BSE ladder. This is quite hard for a non contactlike vertex and, at present, certainly beyond a consistent chiral unitary treatment. Fortunately, available calculations of $\bar{K}N$ scattering within the BSE chiral unitary approach (none of them including box graphs [21–28]) yield a fairly good quantitative description of s-wave baryonic resonance data. This would suggest that the box mechanism would not play a crucial role. In summary, we do not expect the

corrections to the mass and width estimated above for the Θ^* resonance to be large enough to affect its existence. Possible production and identification mechanisms for this resonance could be found in reactions like $\gamma p \to \bar{K}^0 p K^+ \pi^-$ by meassuring the three body $pK^+\pi^-$ invariant mass.

The scheme presented here also contains other exotic states which will be examined elsewhere. For instance in the $Y = -3$, $I = J = 1/2$ sector we find $|\bar{K}^*\Omega\rangle = |\mathbf{1134}; \mathbf{35}_2\rangle$ and thus we have an attractive $\bar{K}^*\Omega$ interaction and possibly a bound state.

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