I review some basic facts about the chiral limit of QCD. This allows to formu-
late an effective field theory below the chiral symmetry breaking scale, chiral per-
turbation theory (CHPT). I show that for threshold reactions, the spectrum of QCD is most eco-
nomically encoded in a set of coupling constants of operators of higher chiral dimension.
A consistent scheme to incorporate the ∆(1232) is also discussed and some examples are given. It is stressed that more precise low–energy data are needed to further test and sharpen the resonance saturation hypothesis in the presence of baryons.

1 Effective field theory of QCD

In the sector of the three light quarks, one can write the QCD Lagrangian as

\[ \mathcal{L}_{\text{QCD}} = \mathcal{L}_{\text{QCD}}^0 - \bar{q} \mathcal{M} q, \tag{1} \]

with \( q^T = (u, d, s) \) and \( \mathcal{M} = \text{diag}(m_u, m_d, m_s) \) the current quark mass matrix. The current quark masses are believed to be small compared to the typical hadronic scale, \( \Lambda_{\chi} \approx 1 \text{ GeV} \). \( \mathcal{L}_{\text{QCD}}^0 \) admits a global chiral symmetry, i.e. one can independently rotate the left– and right–handed components of the quark fields. This symmetry is spontaneously broken down to its vectorial subgroup, \( \text{SU(3)}_{L+R} \), with the appearance of eight massless Goldstone bosons. The explicit chiral symmetry breaking due to the quark mass term gives these particles, identified with the pions, kaons and eta (denoted \( \phi \)), a small mass. The consequences of the spontaneous and the explicit chiral symmetry breaking can be calculated by means of an effective field theory (EFT), called chiral perturbation theory.\( \mathcal{L}_{\text{QCD}} \) is mapped onto an effective Lagrangian with hadronic degrees of freedom,

\[ \mathcal{L}_{\text{QCD}} = \mathcal{L}_{\text{eff}}[U, \partial U, \ldots, \mathcal{M}, B], \tag{2} \]

where the matrix–valued field \( U(x) \) parametrizes the Goldstones, \( \mathcal{M} \) keeps track of the explicit symmetry violation and \( B \) denotes matter fields (like e.g. the baryon octet). While the latter are not directly related to the symmetry breakdown, their

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interactions are severely constrained by the non-linearly realized chiral symmetry and one can thus incorporate them unambiguously. $\mathcal{L}_{\text{eff}}$ admits an energy expansion,

$$\mathcal{L}_{\text{eff}} = \mathcal{L}^{(2)}_{\phi} + \mathcal{L}^{(4)}_{\phi} + \mathcal{L}^{(1)}_{\phi B} + \mathcal{L}^{(2)}_{\phi B} + \mathcal{L}^{(3)}_{\phi B} + \mathcal{L}^{(4)}_{\phi B} + \cdots ,$$

(3)

where the superscript $(i)$ refers to the number of derivatives and/or meson mass insertions. The first two terms in Eq. $(3)$ comprise the meson sector whereas the next four are relevant for processes involving one single baryon. The ellipsis stands for terms with more baryon fields and/or more derivatives. The various terms contributing to a certain process are organized by their *chiral* dimension $D$ (which differs in general from the physical dimension) as follows:

$$D = 2L + 1 \sum_d (d - 2) N^\phi_d + \sum_d (d - 1) N^{\phi B}_d ,$$

(4)

with $L$ the number of (Goldstone boson) loops and $d$ the vertex dimension (derivatives or factors of the pion mass). Lorentz invariance and chiral symmetry demand that $d \geq 2$ ($\geq 1$) for mesonic (pion–baryon) interactions. So to lowest order, one has to deal with tree diagrams ($L = 0$) which is equivalent to the time–honored current algebra (CA). However, we are now in the position of *systematically* calculating the corrections to the CA results. It is also important to point out that $\mathcal{L}^{(4)}_{\phi}$ and $\mathcal{L}^{(2,3,4)}_{\phi B}$ contain parameters not fixed by symmetry, the so–called *low-energy constants* (LECs). These have to be determined from data or are estimated from resonance exchange, as discussed in some detail below. The whole machinery is well documented, see e.g. Refs. 6, 7, 8.

### 2 Decoupling theorem

The chiral limit of QCD exhibits some interesting features. In particular, certain quantities show a non–analytic behaviour in terms of the quark masses. A well-known example is the lowest order Goldstone boson loop contribution to the baryon mass,

$$\delta m_B = \text{const} \cdot M^3_{\phi} \sim m^{3/2}_\text{quark} .$$

(5)

The constant in front of the $M^3_{\phi}$ is given in terms of known parameters. More important, if one would calculate the same one loop diagram with massive (resonance) intermediate states, its value would remain unchanged but some higher order corrections ($\sim M^4_{\phi}$) would get modified. This is the essence of the *decoupling theorem* derived some 15 years ago by Gasser and Zepeda. It states that the *leading* non–analytic corrections (LNAC) to S–matrix elements and transition currents are given solely in terms of $\mathcal{L}_{\text{eff}}[U, B]$ (in the chiral limit), i.e. that the inclusion of mesonic ($\rho, \omega, \ldots$) or baryonic resonances ($\Delta, N^*, \ldots$) in the pertinent loop graphs does not modify the LNACs. Note also that tree graphs are obviously analytic in the quark
masses. A detailed exposition of these results is given in Ref. Another such effect is observed in the electromagnetic polarizabilities of the proton. To lowest order in the chiral expansion, these are given by one loop graphs and thus only depend on known parameters,\(^1^0\)

\[
\begin{align*}
\bar{\alpha}_p &= 10 \bar{\beta}_p = \frac{5e^2 g_A^2}{384\pi^2 F_\pi M_\pi} = \frac{5C g_A^2}{4M_\pi} = 12.4 \cdot 10^{-4} \text{fm}^3,
\end{align*}
\]

which scales like \(1/\sqrt{m_{\text{quark}}}\) and is in good agreement with the empirical numbers, \(\bar{\alpha}_p = (12.1 \pm 0.8 \pm 0.5) \cdot 10^{-4} \text{fm}^3\), \(\bar{\beta}_p = (2.1 \pm 0.8 \pm 0.5) \cdot 10^{-4} \text{fm}^3\). However, it is well known that tree graphs with an intermediate \(\Delta(1232)\) contribute roughly +10 \cdot 10^{-4} \text{fm}^3 to \(\bar{\beta}_p\) and one thus might question the relevance of the CHPT result. Here, the decoupling theorem and pion loops come to ones rescue. As shown in Ref., the large next-to-leading order tree \(\Delta(1232)\) contribution to \(\bar{\beta}_p\) is almost completely cancelled by a pion loop graph at the same order with a subleading non–analyticity \(\sim \ln M_\pi\), which has a large positive coefficient,

\[
\bar{\beta}_p = \frac{C g_A^2}{8M_\pi} + \frac{C}{\pi} \left[ \left( \frac{3(3 + \kappa_s)g_A^2}{m} - c_2 \right) \ln \frac{M_\pi}{\lambda} \right] + \delta\bar{\beta}_p^r(\lambda),
\]

where I have dropped a small finite piece from the loops. The value for \(c_2\) is given below and \(\bar{\beta}_p^r\) subsumes the large \(\Delta\)-contribution from the dimension four counter terms of the type \(\mathcal{L}_{\pi N}^{(4)} \sim N F_{\mu \nu} F^{\mu \nu} N\). Yet another neat example has been discussed by Mallik in connection with terms of the type \(M_\pi^4 \ln M_\pi^2\) contributing to the nucleon mass.\(^1^1\)

3 Resonance saturation

In the meson sector at next-to-leading order, the effective Lagrangian \(\mathcal{L}_\rho^{(4)}\) contains ten LECs, called \(L_i\). These have been determined from data in Ref.\(^2\) (for an update, see e.g. Ref.\(^1^4\)). The actual values of the \(L_i\) can be understood in terms of resonance exchange\(^5\), i.e. the renormalized \(L_i^r(M_\rho)\) are practically saturated by resonance exchange \((S, P, V, A)\). In some few cases, tensor mesons can play a role.\(^1^5\) This is sometimes called chiral duality because part of the excitation spectrum of QCD reveals itself in the values of the LECs. Furthermore, whenever vector and axial resonances can contribute, the \(L_i^r(M_\rho)\) are completely dominated by \(V\) and \(A\) exchange, called chiral VMD.\(^1^6\) As an example, consider the finite (and thus scale–independent) LEC \(L_9\). Its empirical value is \(L_9 = (7.1 \pm 0.3) \cdot 10^{-3}\). The well–known \(\rho\)-meson (VMD) exchange model for the pion form factor, \(F_\pi^V(q^2) = M_\rho^2/(M_\rho^2 - q^2)\) (neglecting the width) leads to \(L_9 = F_\pi^V/(2M_\rho^2) = 7.2 \cdot 10^{-3}\), in good agreement with the empirical value. Even in the symmetry breaking sector related to the quark mass, where only scalar and (non-Goldstone) pseudoscalar mesons can contribute,
resonance exchange helps to understand why SU(3) breaking is generally of $O(25\%)$, except for the Goldstone boson masses. In principle, these LECs are calculable from QCD, a first attempt using the lattice has been reported in Ref. and these studies are continuing.

Matters are much more complicated in the baryon sector, largely due to the fact that one can have baryonic ($N^*$) as well as mesonic ($M$) excitations, symbolically

$$\tilde{L}_{\text{eff}}[U, B; M, N^*] \rightarrow L_{\text{eff}}[U, B],$$

so that the LECs are given in terms of masses ($m_M, m_{N^*}$) and mesonic as well as baryonic coupling constants. Of the baryon resonances, the $\Delta(1232)$ plays a particular role as explained below. Consider first an example where resonance saturation works fairly well, namely the reaction $\gamma p \rightarrow \pi^0 p$ in the threshold region. The most accurate CHPT calculation performed so far contains three LECs, two related to the S-wave $E_{0+}$ (called $a_1$ and $a_2$) and one related to the P-wave $P_3$, called $b_p$. Note that in the threshold region one is only sensitive to the sum $a_1 + a_2$ of the S-wave LECs. A fit to the TAPS data leads to $a_1 + a_2 = 6.60 \text{GeV}^{-4}$ and $b_p = 13.0 \text{GeV}^{-3}$. Note that the SAL data require a somewhat larger value for $b_p$. Resonance exchange leads to a completely fixed vector meson contribution and to one from the $\Delta$, which depends on some off-shell parameters. Constraining these from previous investigations of the proton magnetic polarizability and the $\pi N$ P-wave scattering volume $a_{33}$, one finds $a_1 + a_2 = (a_1 + a_2)^\Delta + (a_1 + a_2)^V = 3.92 + 2.67 = 6.59 \text{GeV}^{-4}$ together with $b_p^{\Delta + V} = 13.0 \text{GeV}^{-3}$ in very good agreement with the numbers obtained in the free fit. Still, there remain uncertainties to be clarified. Consider e.g. the four finite SU(2) LECs $c_{1,2,3,4}$ related to the dimension two pion–nucleon Lagrangian. These have been determined from one loop calculations of the $\sigma$–term and certain $\pi N$ scattering observables to order $q^3$ but also from the subthreshold expansion of the $\pi N$ scattering amplitude to order $q^2$. The resulting values for the $c_i$ are typically a factor 1.5 smaller than in Ref. We have recently reevaluated these coupling constants by constructing observables which to one loop order $q^3$ are given entirely by tree graphs with insertions from the dimension one and two Lagrangian and have finite loop contributions, but are free of insertions from $L_{\pi N}^{(3)}$. The central values are listed in table [1]. These values can also be understood from resonance exchange assuming only that $c_1$ is entirely saturated by scalar meson exchange. Amazingly, the ratio of the scalar meson mass to scalar–meson–nucleon coupling needed, $M_S/\sqrt{g_S} = 180 \text{MeV}$, is exactly the one found in the Bonn potential, where the scalar–isoscalar $\sigma$–meson models the strong pionic correlations in the presence of nucleons ($M_\sigma = 550 \text{MeV}$ and $g_\sigma^2/(4\pi) = 7.1$). Notice also that resonance saturation via vector meson exchange works well for the anomalous magnetic moments. In the isoscalar case, we find $\kappa_s = -0.12$ compared to $\kappa_s = -0.16$. For the isovector moment, the empirical value ($\sim 5.8$) agrees well with the large tensor–to–vector coupling of the $\rho$, $\kappa_\rho \simeq 6$. Also, Mojžiš has recently analyzed other $\pi N$ scattering observables and finds somewhat different
Table 1. Values of the dimension two LECs $c_i' = 2mc_i (i = 1, \ldots, 5)$ as determined in\cite{23} The uncertainties on these parameters are discussed in detail in that reference. The * denotes an input quantity.

| $c_i'$ | occurs in | determined from | central value | res. exch. |
|--------|-----------|----------------|--------------|------------|
| $c_1'$ | $m_N, \sigma_{\pi N}, \gamma N \rightarrow \gamma N$ | phen. + res.exch. | $-1.7$ | $-1.7^*$ |
| $c_2'$ | $\pi N \rightarrow \pi(\pi)N, \gamma N \rightarrow \gamma N$ | phen. + res.exch. | $6.7$ | $7.3$ |
| $c_3'$ | $\pi N \rightarrow \pi(\pi)N, \gamma N \rightarrow \gamma N$ | phen. + res.exch. | $-10.1$ | $-9.8$ |
| $c_4'$ | $\pi N \rightarrow \pi(\pi)N$ | phen. + res.exch. | $6.8$ | $6.6$ |
| $c_5'$ | $(m_n - m_p)^{\text{strong}}, \pi^0 N \rightarrow \pi^0 N$ | phen. | $-0.17$ | $-$ |
| $c_6$ | $\kappa_p, \kappa_n$ | phen. + res.exch. | $5.8$ | $6.1$ |
| $c_7$ | $\kappa_p, \kappa_n$ | phen. + res.exch. | $-3.0$ | $-3.1$ |

This remains to be clarified. In the three–flavor sector, the LECs related to the symmetry breaking terms in the meson–baryon Lagrangian can not yet simply be understood in terms of scalar meson exchange. One would either need unphysically large scalar–baryon couplings or low scalar meson masses (for details, see Ref.\cite{27}). At present, however, one has to use resonance saturation in actual calculations since too few accurate low–energy data exist to pin down all LECs (or at least a subset for $\pi N$ scattering and photo reactions). An example where resonance saturation seems to work even for SU(3) is shown in Fig.\cite{28} In that reference, kaon photo– and electroproduction off protons is considered. A few of the 13 LECs could be fixed from single nucleon properties, the larger number from resonance exchange. Amazingly,

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{recoil_polarization}
\caption{Recoil polarization for $\gamma p \rightarrow K^+\Lambda$ calculated in SU(3) HBCHPT to order $q^3$.}
\end{figure}

the shape of the recoil polarization for $\gamma p \rightarrow K^+\Lambda$ measured at ELSA\cite{29} is well
described (even so the energy at which it was measured is somewhat too high). It is important to stress that for threshold reactions, the resonances need not be taken as dynamical degrees of freedom in the effective field theory but rather should be used to estimate the coefficients of higher dimension operators. However, if one is interested in extending the range of applicability of CHPT say in the region of the \( \Delta(1232) \), it has to be accounted for dynamically.

4 Above threshold: Inclusion of the \( \Delta(1232) \)

Although the inclusion of the decuplet was originally formulated for SU(3),\(^{34}\) let us focus here on the simpler two–flavor case. Among all the resonances, the \( \Delta(1232) \) plays a particular role for essentially \( \text{two} \) reasons. First, the \( N\Delta \) mass splitting is a small number on the chiral scale of 1 GeV,

\[
\Delta \equiv m_\Delta - m_N = 293 \text{ MeV} \simeq 3F_\pi ,
\]

(9)

and second, the couplings of the \( N\Delta \) system to pions and photons are very strong,

\[
g_{N\Delta\pi} \simeq 2g_{NN\pi} .
\]

(10)

So one could consider \( \Delta \) as a small parameter. It is, however, important to stress that in the chiral limit, \( \Delta \) stays finite (like \( F_\pi \) and unlike \( M_\pi \)). Inclusion of the spin–3/2 fields like the \( \Delta(1232) \) is therefore based on phenomenological grounds but also supported by large–\( N_c \) arguments since in that limit a mass degeneracy of the spin–1/2 and spin–3/2 ground state particles appears. Recently, Hemmert, Holstein and Kambor\(^{31}\) proposed a systematic way of including the \( \Delta(1232) \) based on an effective Lagrangian of the type \( \mathcal{L}_{\text{eff}}[U, N, \Delta] \) which has a systematic “small scale expansion” in terms of \( \text{three} \) small parameters (collectively denoted as \( \epsilon \)),

\[
\frac{E_\pi}{\Lambda}, \frac{M_\pi}{\Lambda}, \frac{\Delta}{\Lambda} ,
\]

(11)

with \( \Lambda \in [M_\rho, m_N, 4\pi F_\pi] \). Starting from the relativistic pion–nucleon-\( \Delta \) Lagrangian, one writes the nucleon \( (N) \) and the Rarita–Schwinger \( (\Psi_\mu) \) fields in terms of velocity eigenstates (the nucleon four–momentum is \( p_\mu = mv_\mu + l_\mu \), with \( l_\mu \) a small off–shell momentum, \( v \cdot l \ll m \) and similarly for the \( \Delta(1232) \)\(^{32}\)),

\[
N = e^{-i mv \cdot x} (H_v + h_v) , \quad \Psi_\mu = e^{-i mv \cdot x} (T_{\mu v} + t_{\mu v}) ,
\]

(12)

and integrates out the “small” components \( h_v \) and \( t_{\mu v} \) by means of the path integral formalism developed in Ref.\(^{33}\). The corresponding heavy baryon effective field theory in this formalism does not only have a consistent power counting but also \( 1/m \) suppressed vertices with fixed coefficients that are generated correctly (which is much
simpler than starting directly with the “large” components and fixing these coefficients via reparametrization invariance). Since the spin–3/2 field is heavier than the nucleon, the residual mass difference $\Delta$ remains in the spin–3/2 propagator and one therefore has to expand in powers of it to achieve a consistent chiral power counting. The technical details how to do that, in particular how to separate the spin–1/2 components from the spin–3/2 field, are given in Ref. 31.

5 A few examples

In this paragraph, I briefly discuss two observables which have been calculated in the extension of CHPT including the $\Delta(1232)$.

The first one is related to the scalar nucleon form factor. In fact, this calculation predates the formalism developed by Hemmert et al., i.e. the $\Delta(1232)$ is treated as a heavy spin–3/2 field in the formalism proposed by Jenkins and Manohar. Nevertheless, the pertinent results have been expanded in powers of $\Delta$ and thus the essence of the small scale expansion is captured. Consider the pion–nucleon $\sigma$–term,

$$\sigma_{\pi N}(t) = \frac{1}{2} (m_u + m_d) \langle p' | \bar{u}u + \bar{d}d | p \rangle , \quad t = (p' - p)^2 .$$

(13)

Of particular interest in the analysis of $\sigma_{\pi N}$ is the Cheng–Dashen point, $t = 2M_\pi^2$, $\nu = 0$ (at this unphysical kinematics, higher order corrections in the pion mass are the smallest) and one evaluates the scalar form factor

$$\Delta \sigma_{\pi N} \equiv \sigma_{\pi N}(2M_\pi^2) - \sigma_{\pi N}(0) .$$

(14)

To one loop and order $O(\epsilon^3)$, $\Delta \sigma_{\pi N}$ is free of counter terms and just given by simple one loop diagrams:

$$\Delta \sigma_{\pi N} = \frac{3g^2 M_\pi^2}{64\pi^2 F_\pi^2} \left\{ \pi M_\pi + (\pi - 4)\Delta - 4\sqrt{\Delta^2 - M_\pi^2} \ln \left( \frac{\Delta}{M_\pi} + \sqrt{(\frac{\Delta}{M_\pi})^2 - 1} \right) + \ldots \right\}$$

$$= (7.4 + 7.5) \text{ MeV} \simeq 15 \text{ MeV} ,$$

(15)

where the first term comes from the intermediate nucleon and scales as $M_\pi^3$ whereas the other terms come from the loop graph with the intermediate $\Delta(1232)$ and scale as $\Delta M_\pi^2$, which are therefore both $O(\epsilon^3)$. In the chiral expansion of QCD, however, the first term is $O(q^3)$ while the second is of order $q^4$. This result agrees nicely with the one of the recent dispersion–theoretical analysis (supplemented by chiral symmetry constraints) of Gasser, Leutwyler and Sainio, $\Delta \sigma_{\pi N} = (15 \pm 1) \text{ MeV}$. Notice, however, that the CHPT result might be strongly affected by higher order effects and SU(3) breaking as the $O(q^4)$ analysis of the baryon masses and $\sigma$–terms
presented in Ref.\textsuperscript{27} indicates. For example, the chiral expansion of the $\sigma$–term carried out to second order in the quark masses and for three flavors is\textsuperscript{27}

$$\sigma_{\pi N}(0) = 58.3 \left(1 - 0.56 + 0.33\right) \text{MeV} = 45 \text{MeV}$$  \hspace{1cm} (16)

which shows a moderate convergence but also indicates the importance of the terms of $O(q^4)$. A calculation in the formalism of Hemmert et al. to order to $O(\epsilon^4)$ could help to clarify the situation concerning the higher order corrections.

The second example is the calculation of the electric dipole amplitude $E_{0+}$ for $\pi^0$ production off protons at threshold\textsuperscript{36} This calculation has only been performed to order $\epsilon^3$ so far and can thus not yet compete with the $q^4$ calculations done in the chiral expansion. It nevertheless shows some interesting features of the small scale expansion. The diagrams with intermediate nucleons contributing to this order have been first worked out in Ref.\textsuperscript{37} and later in the heavy baryon approach in Ref.\textsuperscript{33} In the $\epsilon$ expansion, there are a few more graphs, but only the Born graph with an intermediate $\Delta(1232)$ contributes at threshold. The contribution of this graph to $E_{0+}$ scales as

$$E_{0+} \sim V_{N\Delta\gamma} \cdot S_\Delta \cdot V_{N\Delta\pi} \sim \frac{M_\pi^3}{M_\pi + \Delta} \sim M_\pi^2 \sim \epsilon^2$$  \hspace{1cm} (17)

while in the chiral expansion this term would be of order $M_\pi^3$ in $E_{0+}$. The complete result to this order takes the form\textsuperscript{36}

$$E_{0+}^{\text{thr}} = - C \mu \left\{1 - \left[\frac{3 + \kappa_\pi}{2} + \left(\frac{m}{4F_\pi}\right)^2 - \frac{4b_1 \hat{g}_{N\Delta\pi} m}{9g_{\pi N} F_\pi M_\pi + \Delta}\right]\mu + \mathcal{O}(\mu^2)\right\},$$  \hspace{1cm} (18)

with $C = e g_{\pi N}/(8\pi m)$, $\mu = m/M_\pi \sim 1/7$ and $\hat{g}_{N\Delta\pi} = 1.5$. The coupling constant $b_1$ can be extracted from $\Delta N\gamma$ dynamics. Numerically, this new contribution is small, $\delta E_{0+}^{\text{thr}} \sim -0.2 \cdot 10^{-3}/M_\pi$, as it is expected from the analysis where the $\Delta(1232)$ is used to estimate the LECs. However, to this order in the small scale expansion, this expression is complete. It serves as a good example of the previous statement that for threshold observables, it is more economical to use the resonances as frozen degrees of freedom in terms of LECs. The real virtue of including the $\Delta(1232)$ in the effective Lagrangian comes when one considers e.g. photoproduction above threshold.

6 Outlook

Presently, we are working on evaluating the $E2/M1$ ratio in the region of the $\Delta$ pole and the corrections to the $P$–wave LETs for pion photoproduction due to the $\Delta(1232)$ to order $\epsilon^3$.\textsuperscript{38} Most of the work is already done, but results can not yet be reported. It is important to stress how our calculation differs from the one of Butler et al.\textsuperscript{39}, who found a range for $E2/M1$, $4% \leq |E2/M1| \leq 9%$ and an imaginary part
which is of the same size than extracted from the new MAMI data via speed plot techniques. First, they work in SU(3) and include the spin–3/2 decuplet without expanding in $m_{3/2} - m_{1/2}$. Second, one counter term is simply dropped and $1/m$ corrections are not consistently included. Third, the remaining parameters are chosen in ranges such that kaon loop contributions are effectively suppressed. In the light of the work presented in it is certainly necessary to repeat this calculation for SU(2) taking into account all terms to order $\epsilon^3$. Even when that is done, one still might have to go one order further. We hope to be able to report on these results soon.

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