Eur. Phys. J. C (2008) 58: 47–56
DOI 10.1140/epjc/s10052-008-0742-z

THE EUROPEAN
PHYSICAL JOURNAL C

Regular Article - Theoretical Physics

Role of bare propagator poles in phenomenological Dyson–Schwinger type models

S. Ceci[1,2], **A. Švarc**[1,a], **B. Zauner**[1]

[1]Rudjer Bošković Institute, Bijenička c. 54, 10 000, Zagreb, Croatia
[2]*Present address:* Department of Physics and Astronomy, University of Georgia, Athens, GA 30602, USA

Received: 1 April 2008 / Revised: 26 August 2008 / Published online: 2 October 2008

Abstract The generation of dressed meson–nucleon scattering-matrix poles is presented. A possible scenario for the interrelation of bare and dressed baryon poles is shown by using a particular version of coupled-channel Dyson–Schwinger type model. These findings are then applied to the Roper resonance, and the conclusion is drawn that it is dynamic in nature. A possible correlation between bare and dressed propagator poles on one side and the quantities of constituent quark-model calculations on the other side are discussed.

PACS 14.20.Gk · 12.38.-t · 13.75.-n · 25.80.Ek · 13.85.Fb · 14.40.Aq

1 Introduction

Establishing a well defined point of comparison between experimental results and quark-model predictions has for decades been one of the main issues in hadron spectroscopy; yet the present status is still not satisfactory. Experimental results, via partial-wave and amplitude analysis, give reliable information on dressed propagator singularities ("measured" scattering-matrix poles) [1–5], while quark-model calculations usually give information on the three-quark excited states spectrum in the first-order impulse approximation (bare/quenched mass spectrum). It seems obvious that there is no particular reason for these two quantities to be directly compared. Up to now, in the absence of a better recipe, quenched quark-model states have indeed been compared only with the dressed poles' spectrum [6], but the awareness has ripened that a clear distinction between the two has to be made. One either has to perform "un-quenching" (to "dress" quenched quark-model excited states) and compare the outcome to the scattering-matrix poles, or try to take into account all self-energy contributions which are implicitly included in the "measured" scattering-matrix pole parameters, make a model independent undressing procedure and compare the "bare" outcome to the impulse approximation quark-model calculations. The first approach seems to be feasible but highly involved [7]. The latter one, however, seems to be impossible [8] due to very general field-theory considerations [9–15]. The present recommendation given at the BRAG2007 workshop [16] is that

> "...dressed scattering-matrix singularities are the best, model independent meeting point between quark-model predictions and experiments, and bare quantities in coupled-channel models remain as legitimate quantities to be extracted only within a framework of a well defined model. To understand and interpret them correctly, one has to keep track of the existence of the hadronic mass shifts produced by off-shell-ambiguities, and take them fully into account...".

In this paper, we focus on exploring the detailed properties of both, dressed and bare scattering-matrix parameters in coupled-channel models, and investigate their layout and interdependence. The main purpose of this article is twofold: (i) to study the causal connection between bare and dressed poles and correlate it with the phenomenon of dynamically generated resonant states [17–22], and (ii) to analyze the spectrum of bare propagator poles and to make a connection between bare propagator poles and quark-model states. The design of coupled-channel Dyson–Schwinger models is very convenient for tackling both of these issues, so we shall use it throughout this paper.

Within all coupled-channel formalisms dressed and bare propagator poles are automatically distinguished, but still directly connected. The same type of Dyson–Schwinger equation is always solved, but the treatment of a channel-resonance vertex interaction varies from phenomenological

[a] e-mail: alfred.svarc@irb.hr

to microscopic [23]. The resolvent equation formalism generates self-energy correction terms; self-energy terms shift the initially real-value bare propagator poles into the complex energy plane; and eventually, dressed propagator poles are generated. The obtained full scattering-matrix poles are (at least in principle) experimentally identifiable, but it is clear that the whole chain of self-energy contributions lies between them and the basic microscopic interaction structures.

A direct interrelation of bare and dressed propagator poles offers new insight into the genesis of dressed resonant states. In the coupled-channel formalisms it turns out that dressed scattering-matrix poles can be generated in two very different scenarios. In both cases they are produced when the bare propagator poles are shifted from the real axes into the complex energy plane by a self-energy term produced by an intermediate state excitation of all virtual two-body states. In the first case, the self-energy term is small, so we can readily identify the actual, directly responsible nearby bare propagator pole. In the second case, the self-energy term is large, and a nearby bare term responsible for the dressed pole creation cannot be identified. This offers a simple mechanism to distinguish between *genuine resonances*, defined as states produced by a nearby bare propagator pole, and *dynamic* ones for which such a nearby bare pole does not exist.

As a result, we got a better insight into an unresolved issue which has for quite some time been troubling the physics community: understanding the nature of the Roper resonance.

There is a prevailing belief that the Roper resonance differs significantly from other resonances (most recently recapitulated in [24]), but a fully convincing argument that it is indeed so has never been given. Possible explanations of the phenomenon have been ranging from understanding the Roper resonance as being a hybrid state with excited glue [25–27], to understanding it as a five-quark state, $qqqq\bar{q}$, which produces a scattering-matrix resonant behavior without a standard three-quark pole term [21] (dynamic resonance). The Roper resonance has strongly attracted our attention in the coupled-channel approach, so we have made a special effort to analyze it in the context of the role of bare propagator poles.

As we have said before, non-trivial analytical properties of the self-energy function result in two kinds of dressed propagator poles: those in relative proximity of bare masses, and those that are far from them. It came as a surprise to see that all but one of the dressed poles located in the N^* energy range are of the first kind, and all but one poles of the second kind are lying far outside the resonant region (region of model applicability), having no influence on the scattering observables whatsoever. And the remaining singular case is the Roper resonance. Hence, the mechanism defined before to distinguish between genuine and dynamic dressed scattering-matrix singularities offers us the opportunity to characterize the Roper resonance as being dynamic in nature.

All efforts to understand the relation of bare propagator poles and quark-model quantities have created an ongoing controversy. Some colleagues attempt to relate them directly. In [28–31] the possibility has been opened that $\gamma N \to \Delta$ helicity amplitudes and transition form factors of the constituent quark models should be compared to the bare coupled-channel functions, and in [32] a simple well defined model is devised for understanding the Roper and Δ resonances in terms of cloudy bag model form factors [33], and in [34, 35] this idea is used for understanding charmed/strange resonant states in the meson–meson scattering sector. Other, more cautious ones give them in a certain measure physical importance, but strongly refrain from giving them such a tempting physical meaning. The main reason for such a disagreement lies in the fact that the poles of the interaction potential do arise only from the model assumed, and as such do not reveal much dynamics of the interaction [8, 16, 36]. However, the possibility that within a well defined model a link might be established between the bare poles and quark-model resonant states is still quite appealing. In order to establish some connection between theory and experiment, we pursue the hypothesis that the bare poles are, if not equivalent, then yet much closer to the quark-model structures than the directly measurable quantities. They should, at least, reflect the quark-model pattern of behavior.

We try to relate bare propagator poles and parameters of the quark model of [6]. However, as the parameters in that publication were adjusted according to baryon spectroscopy data, some quark-model resonant states of that model already incorporate loop effects to a certain, but yet unclear extent. Nevertheless, since no channel-resonance coupling (responsible for the dressing effect) is present in that model, we assume that the quark-model states of [6] should still correspond to our bare states.

At this point let us observe that the structure of the Carnegie–Melon–Berkeley (CMB) type model [37] used by the Zagreb group [38, 39] is very similar to the structure of the dynamical coupled-channel model of [28, 29] where a tempting direct interrelation between bare parameters and quark-model states was first proposed. As the major difference between the two (the intricate Hamiltonian-based description of the reaction mechanism in [28, 29, 31] is in [38, 39] replaced with a phenomenological representation) by no means affects the considerations regarding the interpretation of bare quantities. In this paper we use the CMB type model in order to investigate the aspects discussed before of the bare and dressed propagator poles.

The result of our three-channel CMB type model fitted to πN elastic and $\pi N \to \eta N$ partial-wave data concur with the conclusions that the Roper resonance is a dynamic state.

We also, for several lowest partial waves, find a reasonable correspondence between the bare propagator pole parameters and the quark-model states of [6], as well as between the dressed propagator pole parameters and the experimental N^* parameters of [40].

2 Formalism

We shall not repeat the details of the technically involved formalism given extensively in [38, 39], but, for the convenience of the reader, we shall recapitulate the essence.

We use the Carnegie–Melon–Berkeley (CMB) fully analytic, manifestly unitary multi-channel approach of [37]. The model has two main ingredients: the bare resonant propagator $G_0(s)$—the diagonal real matrix in *resonant* indices incorporating real first-order poles and the channel-propagator $\Phi(s)$—the diagonal complex matrix in *channel* indices with matrix elements $\phi(s)$, which takes care of other non-pole singularities. The solution of the problem, the dressed resonant propagator $G(s)$, for each partial wave contains a mixture of all resonant and background contributions. Resonant contributions are generated by "dressing" bare propagator poles from the N^* energy domain with self-energy terms. One attractive and one (occasionally two) repulsive subthreshold terms coupled to the elastic channel have been used to simulate extra contributions to the left-hand cut and background terms. One (sometimes two) additional "background" resonances which had their s_i values fixed at arbitrary chosen large values above the data region have as well been allowed. The dressed resonant propagator $G(s)$ is obtained by explicitly solving the Dyson–Schwinger equation $G(s) = G_0(s) - G_0(s)\ \Sigma(s)\ G(s)$. The graphical illustration of the equation is given in Fig. 1a. The self-energy term $\Sigma(s)$, as shown in Fig. 1b, is effectively represented with the chosen analytic channel propagators and channel-resonance coupling matrix γ as $\Sigma(s) = \gamma^T \Phi(s)\gamma$, and the model manifestly satisfies unitarity. The unitary-normalized partial-wave T matrix is obtained from the dressed propagator matrix G and (to keep the S matrix unitary) the square root of the real matrix function $\mathrm{Im}\,\Phi$ as $T(s) = \sqrt{\mathrm{Im}\,\Phi(s)}\gamma G(s)\gamma^T \sqrt{\mathrm{Im}\,\Phi(s)}$.

The model contains N_C channels, N_P bare propagator poles s_0 (N_R resonant and two background ones), and all channel-to-resonance mixing matrix real parameters γ. Once the number of channels N_C, and the number of bare propagator poles N_P are chosen, the model contains a total of $N_P + (N_P \times N_C)$ free parameters per partial wave. All parameters of the model, the non-square parameter matrix γ and the values of the real bare propagator poles s_0, are concurrently obtained from the least-square fit of the CMB model $T(s)$ to the experimentally obtained partial-wave data.

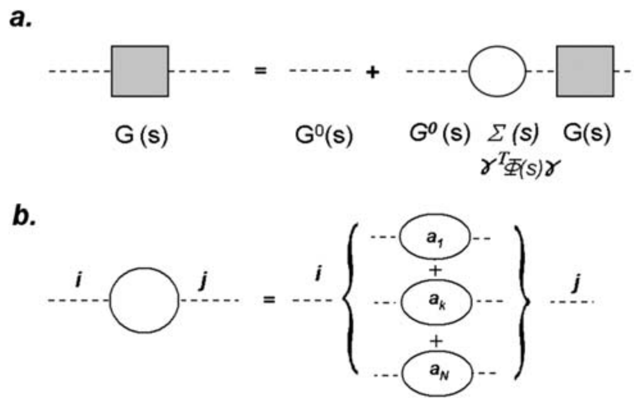

Fig. 1 Resolvent Dyson–Schwinger equation and self-energy term

When fitting, we start with a minimal set of poles: $N_R = 1$ resonant ones, and two (one attractive, one repulsive) subthreshold ones for the background, with the exception of the P_{11} partial wave where we have added the second (repulsive) subthreshold term representing the nucleon pole. Then we increase the number of above-threshold poles until a satisfactory fit is achieved, i.e. until the quality of the fit, measured by the lowest reduced χ^2 value, cannot be further reduced. In addition, visual resemblance of the fitting curve to the data set as a whole is used as a rule of thumb: we reject all those solutions which have a tendency to accommodate for the rapidly varying data points regardless of the χ^2 value. Unsurprisingly, we usually do not obtain a unique solution due to the poor quality of the available data set. However, some poles (usually the lower ones) are quite stable, while the higher ones, exactly where the data are scarce and the need is pronounced for having much more inelastic channels at one's disposal, are "traveling around", producing optically similar partial waves in the fitted channels. We have also allowed for background poles lying high above the data-region energy (we did not restrict the above-threshold pole energy range), but our criteria (lowest reduced χ^2 value, visual resemblance) did not require these terms. We are confident that the better stability will be achieved by improving the quality of the data in presently measured channels, but much more so by increasing the number of channels measured.

So we can safely say that even at this level the singularities of the finally obtained T matrix, the dressed scattering-matrix poles, are the quantities which are constrained directly by experiment.

3 Analysis

To investigate the link between bare and dressed poles, it is important to accentuate three very important features of the CMB type approach:

(1) The causal connection between bare and dressed scattering-matrix singularities is directly built into the formalism
(2) Each dressed scattering-matrix pole is produced by only one bare propagator pole, and
(3) This connection is completely determined by the values of the non-vanishing coefficients of the channel-resonance mixing matrix γ obtained directly as a result of the fit

It is technically very simple to make operational the desired correspondence between bare and dressed scattering-matrix singularities. First we fit the data and obtain the set of bare poles and the parameter values for the channel-resonance mixing matrix γ. Then we establish the existence and location of dressed scattering-matrix poles using one of the proper pole-search methods [41, 42]. Since the Dyson–Schwinger relation in the CMB model mixes all resonances in a partial wave making the self-energy matrix non-diagonal and energy dependent, it seems logical to assume that a simple relation between dressed and bare poles is lost. However, the anticipated and desired link is conveniently established by gradually reducing the values of the mixing parameter matrix from their realistic values to zero by taking the limit $\gamma \to 0$ (we are gradually sliding into the "world without interaction"). The self-energy contribution eventually vanishes, and the values we reach uniquely define which bare poles are responsible for creating the particular dressed ones when the γ matrix is "switched back on" (going back to the "world with interaction").

Consequently, it is straightforward to see which dressed pole is produced by a particular bare one *merely by numerically decreasing the values of the γ matrix mixing parameters from their realistic values obtained in the fit to zero.*

This causal correspondence between dressed and bare poles is a mathematical necessity, but it is *not at all* self-evident that

(1) *All physically observable* dressed singularities *must* be created by nearby bare ones, and neither is it clear that
(2) *All* bare singularities *must* remain near the physical axes and create *physically observable* dressed ones

In our model we have found counter-examples for both.

As a corollary of (1), our model offers us a possibility to define two types of dressed scattering-matrix poles: (a) the *genuine* resonant state as a state which is produced by a nearby bare propagator pole, and (b) the *dynamic* resonant state as a state which is created out of a distant bare propagator pole through the interaction mechanism. The naming *dynamic resonance* is used in analogy to [21].

As a corollary of (2) we perceive that some of the bare propagator poles might de facto "get lost". Namely, the resonance mixing, as well as other self-energy features might shift the bare poles from the real axis far enough into the unphysical range so that they lose their possible physical interpretation. In order to shift a bare pole far enough in the complex plane so that we may find it missing, the dressing has to be fairly strong. The poles are either shifted far away into the complex direction so that their influence upon the T matrices on the real axes is smooth, or they are pushed outside the analyzed energy range. In both cases they evade identification. The easiest way to achieve this would be to make the coupling constants unrealistically large. Nevertheless, as we shall show later, in the case of $N(1440)$, there are other possibilities as well. The way in which the self-energy is created (matrix multiplication of resonant and background terms) as well as their non-linear and non-trivial energy dependence drastically change the position of dressed T matrix poles producing fairly large shifts. Moreover, as shown in [43], even substantially large coupling constants do not have to be considered entirely unrealistic. Therefore, the number of bare poles may be different from the number of dressed ones in the measurable energy region (i.e. experimentally observed resonances). Hence, if we manage to link bare poles and quark-model states, we have found another mechanism how to understand the "missing resonance" problem.

Possible interpretation of bare propagator pole spectrum in terms of quenched quark-model states is altogether another affair.

Similarly to the approach presented in [28, 29, 31], we are also unable to prove that bare propagator parameters directly correspond to quenched quark-model states. Such a task would demand a full understanding of how quark models build up the meson–baryon interaction on the hadronic level, and this is unfortunately still beyond our capabilities. For now, we put forth the hypothesis that some correlation between the two should exist, and we try to see where it might take us. We relate real valued bare poles and quenched quark-model states, and imaginary parts of the dressed poles (produced when a complex self-energy $\Sigma(s)$ pushes quark-model states into the complex energy plane) and the decay width. Let us observe that, in addition to creating the imaginary part, bare propagator poles (quark-model state masses) are shifted as well.

At this moment, we feel compelled to emphasize that putting forth such a hypothesis calls for a substantial amount of caution, because not all things are going that favorably. As has recently been pointed out [8, 16], the self-energy term, containing off-shell effects, seems to be explicitly model dependent. It is very likely that the well-known invariance of the scattering matrix with respect to the field transformation of the effective Lagrangian [10–15] makes off-shell effects and the contact terms mutually inter-transformable at the level of the same power counting diagrams, and the self-energy term becomes defined up to a model dependent hadronic shift only. Its absolute value, hence, becomes strictly linked to the choice of the model, and the uniqueness of separating dressed pole into the bare and self-energy contributions disappears. Consequently, it is to be expected that the absolute scale must cease to be well defined. Bearing that in mind, we are strongly interested to see whether, if absolute values of bare poles and quark-model states are not directly related, then at least their grouping remains preserved. To achieve this, we primarily look for the correlations in the number of states (bare poles versus quark-model states) and their relative separation. Absolute positions are to be taken with reserve.

By accepting the hypothesis of linked bare poles and quenched quark-model states quite a number of issues is put into order: dressed poles are through bare poles directly linked to the quark-model states; the nearest quark-model state responsible for a specific scattering-matrix pole can be identified; it is seen how the interaction can push the "reasonable" quark-model state into the complex energy domain *not accessible* to present experiments (the state "gets lost"); the mechanism is revealed by which the scattering-matrix poles can be created not from quark-model states, but otherwise.

As we are dealing with effective interaction theories where all unknowns are just represented by the free parameters of the model, we do not expect the identification to be ideal. The incompleteness of the channel basis is an important issue that will reduce the predictability of the CMB model. What we hope is to get a glimpse of the mechanism how the dressing mechanism shifts the initial pole values, and to get a feeling for its significance. We wonder how large the shift actually is, and if it is present at all.

Furthermore, a sensitive question is the stability of the Zagreb CMB model. In this paper the notion of stability can only be discussed in the context of the stability of the resonant-parameter extraction from the chosen particular partial-wave data set, and not at all in the sense of a stability of the result with respect to a possible change of input. For the Zagreb CMB model, being based on the BATINIC 95 formalism which is accepted and quoted as one of the recognized results in all PDG-RPP issues since 1998, such a stability is (similarly to all other partial-wave analyses in PDG [40]) better for lower (lighter) resonances than for higher (heavier) ones. As the role of the input data set for the stability of the solution is a complex problem right now under study by the Zagreb group, we hope to present the whole discussion in a more coherent way in a forthcoming publication.

4 Results and discussion

We use a model with three channels: two physical two-body channels πN and ηN, while the third, effective channel represents all remaining two- and three-body processes in a form of a two-body process.

For the πN elastic partial waves we used the VPI/GWU single-energy solutions [45, 46].

For the $\pi N \to \eta N$ partial-wave data we used the coupled-channel amplitudes from Batinić et al. [38, 39], but instead of using smooth theoretical curves, we constructed the data points by normally distributing the model input (as in [47]).

The fitting strategy was taken from [47].

The obtained curves correctly reproduce all input partial-wave data and are given in Figs. 2 and 3 for the πN elastic and $\pi N \to \eta N$ process.

The first four partial waves in the $I = 1/2$ channel (S_{11}, P_{11}, P_{13} and D_{13}) were analyzed.

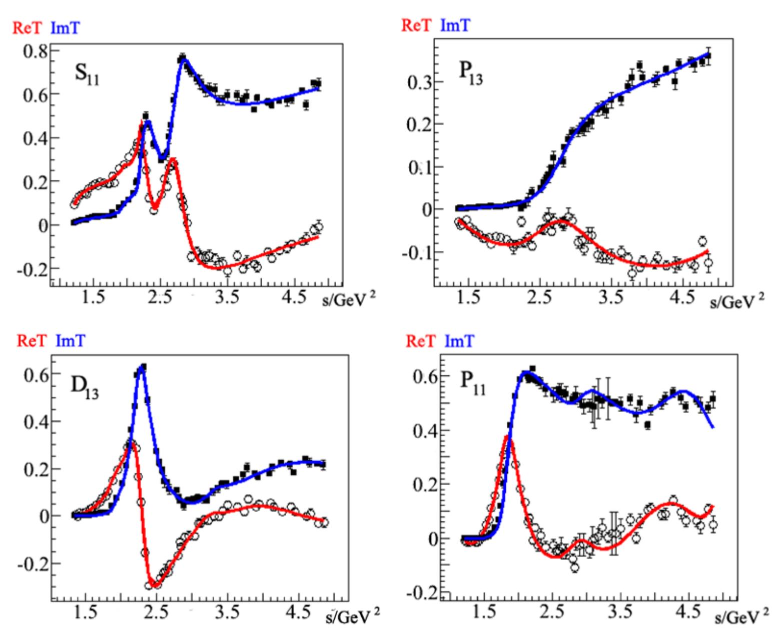

Fig. 2 The obtained curves for πN elastic scattering. *Empty squares* and *red lines* denote the real part of the scattering matrix and *full squares* and *blue lines* denote the imaginary one

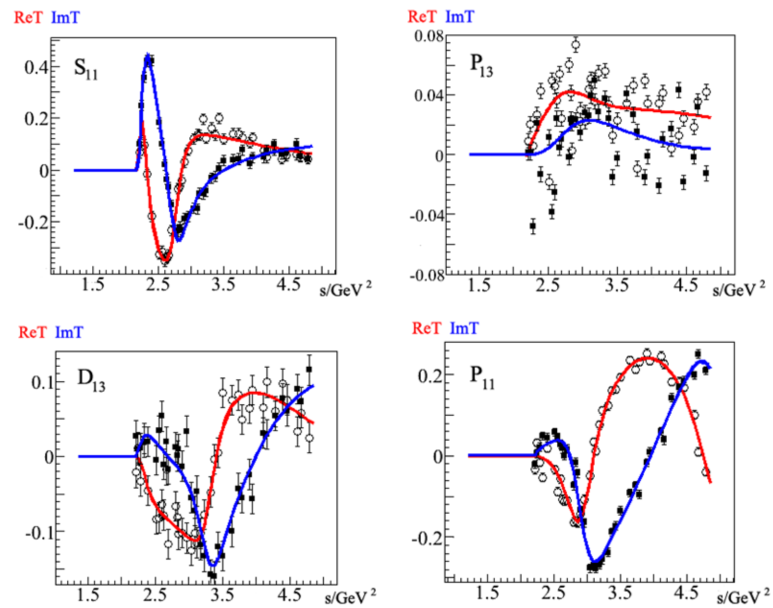

Fig. 3 The curves obtained for the $\pi N \to \eta N$ process. The notation is the same as in Fig. 2

Table 1 Bare propagator and scattering-matrix poles. Stable poles are given in boldface, and non-stable ones in grey. An asterisk in P_{13} value from [40] indicates that the Breit–Wigner mass and width are given instead of the pole position

Partial wave	Bare propagator poles $w = \sqrt{s_0}$ *quark model masses* [48] (GeV)				Scattering-matrix poles—our model *scattering-matrix poles* [40] (GeV)				
	No. 1 (w)	No. 2 (w)	No. 3 (w)	No. 4 . . . (w)	No. 1 $\binom{\mathrm{Re}\,w}{-2\,\mathrm{Im}\,w}$	No. 2 $\binom{\mathrm{Re}\,w}{-2\,\mathrm{Im}\,w}$	No. 3 $\binom{\mathrm{Re}\,w}{-2\,\mathrm{Im}\,w}$	No. 4 $\binom{\mathrm{Re}\,w}{-2\,\mathrm{Im}\,w}$	No. 5 $\binom{\mathrm{Re}\,w}{-2\,\mathrm{Im}\,w}$
S_{11}^{-}	**1.518**	**1.651**	1.912	–	$\boldsymbol{\binom{1.51}{0.10}}$	$\boldsymbol{\binom{1.65}{0.14}}$	$\binom{1.91}{0.95}$	–	–
	1.460	*1.535*	*1.945*	*2.030, . . .*	$\binom{1.51}{0.10}$	$\binom{1.65}{0.14}$	$\binom{2.15}{0.35}$	–	–
D_{13}^{-}	1.477	**1.663**	1.934	–	$\binom{1.40}{0.15}$	$\boldsymbol{\binom{1.51}{0.11}}$	$\binom{1.82}{0.16}$	–	–
	1.495	*1.625*	*1.960*	*2.055, . . .*	$\binom{1.51}{0.11}$	$\binom{1.68}{0.10}$	$\binom{1.88-2.05}{0.16-0.20}$	–	–
P_{13}^{+}	**1.909**	2.484	–	–	$\boldsymbol{\binom{1.65}{0.36}}$	$\binom{2.47}{0.18}$	–	–	–
	1.795	*1.870*	*1.910*	*1.950, . . .*	$\binom{1.66-1.69}{0.12-0.28}$	$\binom{1.90}{0.50}^{*}$	–	–	–
P_{11}^{+}	0. 960	**1.854**	2.018	2.759	$\binom{1.10}{0.00}$	$\boldsymbol{\binom{1.35}{0.16}}$	$\boldsymbol{\binom{1.70}{0.10}}$	$\binom{2.00}{0.60}$	–
	0. 960	*1.540*	*1.770*	*1.880, . . .*	–	$\binom{1.37}{0.19}$	$\binom{1.72}{0.23}$	$\binom{2.12}{0.24}$	–

The obtained bare propagator and scattering-matrix poles are collected in Table 1.

However, even before opening the discussion, it is important to warn the reader that the simplicity of the model, i.e. the fact that we are using only three out of at least seven open, and potentially important channels, will produce only qualitative results. The complexity of the coupled-channel model (simultaneous mixing of all channels) requires considerably larger number of fitting parameters. The absence of constraining data in more than only a few channels will necessarily produce incomplete solutions. However, including more channels unconstrained by the data would, in this model, produce numerical instabilities in the fitting results obtained. To avoid such an undesirable scenario, it is important to get hold of as many constraining data, in as many channels, as possible [44]. The quality of the data is of importance, but the abundance of the constraining channels is what counts.

4.1 Correlation of bare and dressed propagator poles

The first goal of this research, to see the influence of the interaction onto the bare propagator poles, and their "journey" from the initial bare to the final dressed positions, is symbolically visualized in Figs. 4 and 5. In the world without interaction the γ matrices vanish, and the scattering-matrix poles get "undressed", and become equal to bare poles. In the real world, the γ matrices are non-vanishing, and they are obtained by fitting the partial-wave data. Arrows illustrate the

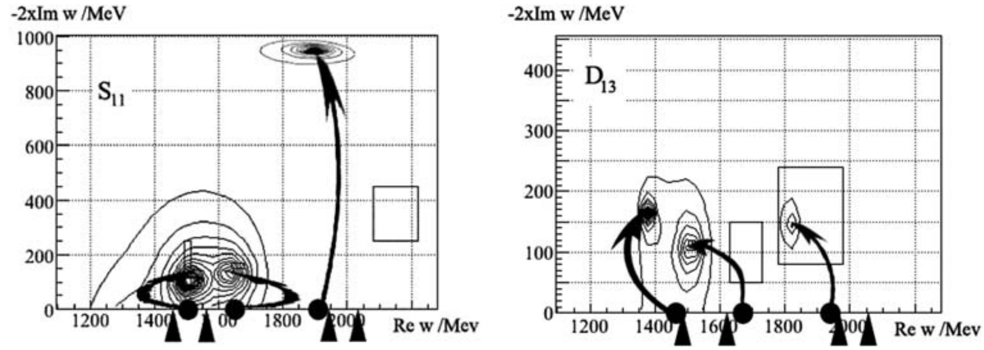

Fig. 4 Scattering-matrix singularities and bare propagator pole positions for the two lowest negative-parity states. *Full dots* denote bare propagator pole positions, *triangle arrows* denote the few lowest quark-model resonant state masses of [48]. *Boxes* represent RPP estimates for dressed pole positions [40]

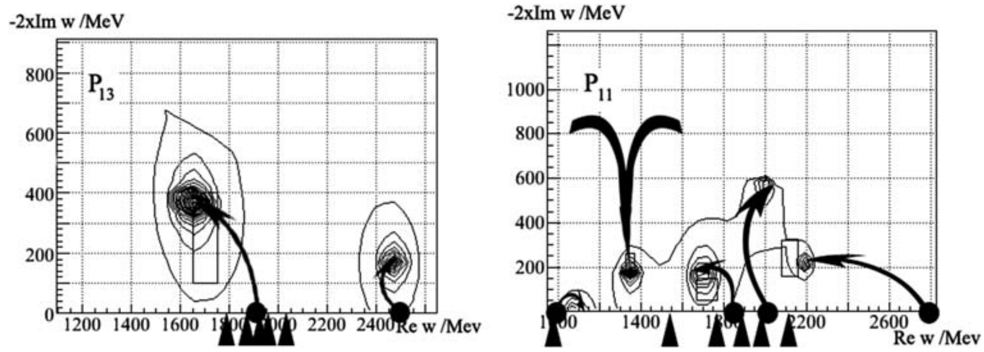

Fig. 5 Scattering-matrix singularities and bare propagator pole positions for the two lowest positive-parity states

way how bare poles travel from the real axis ($\gamma = 0$) towards their final positions in the complex energy plane ($\gamma \neq 0$).

A very interesting conclusion emerges when looking at the results given in Figs. 4 and 5.

All *but one* identified resonant states for the four lowest partial waves seem to be genuine, i.e. there always exists a well defined bare pole which is shifted by the interaction from the real axis into the complex dressed pole. The exception is $N(1440)$ in the P_{11} partial wave, the (in)famous Roper resonance. It is generated differently—as a collective effect of intermediate-state excitations of all virtual two-body states which brings a distant bare background pole back into the N^* energy range. Hence, concerning the often expressed belief that the Roper resonance differs significantly from other resonances, we agree with the view that it might not be a genuine nucleon resonance at all [21]. In the model used here (a CMB-like multi-resonant model with three channels) the Roper resonance turns out to be a *dynamic* resonance.

The advantage of the model we use with respect to the approach of Krehl et al. [21] is that we do not have to go through a cumbersome procedure of adding an extra pole term to the full K matrix coupled-channel effective Lagrange model and demonstrate its superfluousness. Our straightforward criteria—the presence or absence of a nearby bare propagator pole—gives us immediately the decisive answer. Similarly to [21], we in our model do find a scattering-matrix pole in the correct energy regime, but that singularity is not produced by a nearby bare propagator pole, but by a distant background one. In other words, we have no need for a quark-model state corresponding to a Roper resonance.

Another possible explanation why we, in this publication, find a deficit of bare poles with respect to quark-model states, is that only pion- and eta-nucleon channels are considered. It is not at all unreasonable to assume that some quark-model states couple to channels other than these two. Consequently, they should be seen elsewhere [44]. This statement is strongly supported by the fact that a number of new resonances is recently reported when analyzing other channels (the photo-production channel in particular). Let us only mention that, in recent years, quite a number of authors have put forward the need for new resonances in S_{11}, D_{13}, and D_{15} partial waves between 1.73 and 2.1 GeV [49–57], and that their existence cannot be established when using our limited input data set.

4.2 Spectrum of bare propagator poles

The tempting idea has been introduced by Sato and Lee [28, 29, 31] that there exists a direct correspondence between bare poles and quenched quark-model states.

We accept this idea with a grain of salt because of the potentially unpleasant model dependence issue. Bare parameters (bare resonant masses, in this particular case) are inherently model dependent, and this naturally led many authors to draw the simple and natural conclusion: there is no need to investigate anything related to bare propagators

[8, 16]. Nevertheless, we have decided to further investigate the correlation between the spectrum of bare poles and quark-model states in the Zagreb coupled-channel model. However, taking the warnings seriously into account, we do not look for the direct identification of bare propagator poles and quark-model states, we only look for the same *pattern of behavior* of both. We study the *absolute number of each*, and their *relative distance*. The absolute scale is taken with utter reserve.

Regarding the distribution of the bare-pole positions, we present our results for the four lowest partial waves bearing in mind the afore-mentioned ambiguity of the hadronic shift in macroscopic models. Two groups of results are given in Table 1: two lowest negative- and two lowest positive-parity partial waves. The first column contains bare propagator poles produced by the Zagreb CMB model (normal text) and one out of several available sets of quark-model masses (italics) [48]. The second column contains scattering-matrix poles produced by the Zagreb CMB model from this paper (normal text) and their standard values from PDG (italics) [40].

First group of results for the two lowest negative-parity partial waves S_{11} and D_{13} show a reasonable amount of agreement. The number of bare poles needed by the input data, and the number of lowest lying quark-model states from [48] is identical for both partial waves. As can be seen in Table 1 and in Fig. 4, all three bare poles for both of them can be naturally identified with the lowest three-quark states of [48]. There are some discrepancies in the mass position, but each required bare propagator pole does qualitatively correspond to a particular three-quark state, and all lowest ones have found their bare propagator counterpartners.

The obtained scattering-matrix pole positions correspond reasonably well to the "experimental" values reported in [40]. The only disagreement, the unexpected position of the third dressed pole of the S_{11} partial wave (too far in the complex energy plane), is most probably a consequence of the fact that we fit the data from only two processes. It is expected to disappear on including additional channels. All three experimentally detected dressed poles for the D_{13} partial wave are reproduced, but only the second one, at 1.51 GeV, is stable. The used data seem to be insufficient, so either the lowest two states are somewhat shifted in mass, or we predict an extra, yet unseen state at 1.4 GeV, and do not see an experimentally confirmed state at 1.68 GeV.

The next group of results is for the two lowest positive-parity partial waves: P_{11} and P_{13}. The P_{13} wave is not inconsistent with the hypothesis of the article, but quite some problems occur in the P_{11} case.

In the case of P_{13} partial wave, as can be seen in Table 1 and in Fig. 5, only one out of the five three-quark states of [48] is identified with the bare propagator pole, while the other states remain uncorrelated. The shape of the P_{13} wave looks relatively simple, so fits usually do not demand for additional resonances (poles). As discussed before, the positions of higher poles tend to be not too precisely defined due to the incompleteness of the data set, so the second bare propagator pole should either be identified with one of the higher lying quark-model states, or its position could be shifted downwards after including more channels. Of course, such a disagreement might also mean something completely different, like, for instance, it may indicate that our analysis disproves some quark models. However, fitting more channels concurrently should remedy this, since many predicted (yet unobserved) three-quark states are expected to couple much more strongly to channels presently omitted (the famous missing resonance problem).

The notoriously problematic P_{11} partial wave, however, remains troublesome as in the majority of theoretical considerations. The important difference is the existence of a well defined subthreshold pole—the nucleon pole. Following the hypothesis of this paper that we should try to identify bare propagator poles with *all* quark-model states (resonant *and* bound), we have added an extra subthreshold pole to the standard non-resonant background, and fixed its value at 0.96 GeV (the mass of the quark-model subthreshold nucleon pole of [48]). Then, we have left the remaining three poles unconstrained. As shown in Table 1, all known above-threshold dressed poles [40] are reproduced. The first concern is that the dressing procedure shifted bare "nucleon" pole from 0.96 to 1.1 GeV, nowhere near the physical nucleon pole at 0.939 GeV. The fact that we do not obtain the correct value for the nucleon pole might seem to be a problem. However, this we are not attributing things to the deficiency of the model, but rather to the potential problems with the chosen input data set for the P_{11} partial wave (VPI/SAID single-energy solutions for the πN elastic channel). Just to give a clue as to what might be happening: our original solution [38, 39], which has used KH80 for describing the πN elastic channel, has one bare background pole at 1.04 GeV, but a dressed nucleon pole turns out to be at the *correct* value of 0.94 GeV. Unfortunately, these values have still not been published, as their importance has until recently evaded us entirely.

Real problems began when the identification of quark-model resonant states with bare propagator pole positions was attempted. In [47] we have demonstrated that the presence of inelastic channels directly produces the $N(1710)$ P_{11} pole, and in Fig. 5 we show that it is generated by dressing the 1.854 GeV bare propagator pole. This pole can be directly associated with one of the quark-model states of [48], either 1.770 or 1.880. The nucleon state is producing an insignificant, subthreshold and experimentally inaccessible pole at 1.100 GeV. The remaining two bare poles at 2.018 and 2.759 produce the dressed pole at 2.200 GeV which can be identified with the poorly determined $N(2100)$ P_{11}, and an experimentally not yet established state around 2 GeV.

The model used, with constraining data in only two channels, shows two very interesting features for the P_{11} partial wave:

(1) One of the experimentally confirmed dressed poles, namely the $N(1440)$ P_{11} state—a Roper resonance—is not produced by any nearby bare pole as was the case for all other ones.
(2) No bare propagator pole which would correspond to the 1.540 GeV quark-model state is needed.

5 Conclusions

The correlation between bare and dressed poles, and the idea that the quenched quark-model resonant states are to be identified with a set of bare poles obtained in a fit (quark-model states ⇔ bare propagator poles) [28, 29, 31] is tested on a three-channel CMB type model.

For the first few partial waves we show the following.

The $N(1440)$ P_{11} resonance (Roper resonance) is a scattering-matrix pole *not* produced by a nearby bare propagator pole. It is a dynamic resonant state, created as a collective effect of intermediate state excitations of all virtual two-body states which brings a distant bare pole back into the N^* energy range. Thus, it does not seem to be a genuine excited state.

The bare propagator pole distribution nicely corresponds to the spectrum of low-lying negative-parity quark-model states.

The number of bare propagator poles needed to explain the data is, for the low-lying positive-parity states, much lower than the number of corresponding quenched quark-model resonant states. The bare propagator poles needed to explain the "missing resonant states" are to be looked for by including other channels.

The $N(1710)$ P_{11} resonance is a genuine resonance produced by a nearby bare propagator pole which can be identified with a quark-model state. However, its final positioning awaits new data in other inelastic channels.

The bare propagator pole corresponding to the lowest P_{11} quark-model state of [48] could not be identified when only πN elastic and $\pi N \to \eta N$ experimental data are used, so we wonder what data and from what processes ($\pi N \to \pi\pi N$?) could confirm its existence.

References

1. J.S. Ball, R.R. Campbell, P.S. Lee, G.L. Shaw, Phys. Rev. Lett. **28**, 1143 (1972)
2. A. Nagova, J. Pisut, Nucl. Phys. B **61**, 445 (1973)
3. A. Nagova, J. Pisut, Nucl. Phys. B **65**, 544E (1973)
4. T.D. Spearman, Phys. Rev. D **10**, 1660 (1974)
5. J.S. Ball, R.L. Goble, Phys. Rev. D **11**, 1171 (1975)
6. S. Capstick, W. Roberts, Prog. Part. Nucl. Phys. **45**, S241–S331 (2000)
7. S. Capstick, Unquenching the Quark Model. Presentation given on BRAG 2007, September 4, 2007, University of Bonn, Germany. http://www.kph.uni-mainz.de/MAID/BRAG2007/contributions/
8. C. Hanhart, What we can learn from spectroscopy and what not. Presentation given on BRAG 2007, September 4, 2007, University of Bonn, Germany. http://www.kph.uni-mainz.de/MAID/BRAG2007/contributions/
9. H.W. Fearing, S. Scherer, Phys. Rev. C **62**, 034003 (2000)
10. M. Gell-Mann, M.L. Goldberger, Phys. Rev. **96**, 1433 (1954)
11. R. Haag, Phys. Rev. **112**, 669 (1958)
12. J.S.R. Chisholm, Nucl. Phys. **26**, 469 (1961)
13. S. Kamefuchi, L. O'Raifeartaigh, A. Salam, Nucl. Phys. **28**, 529 (1961)
14. S. Coleman, J. Wess, B. Zumino, Phys. Rev. **177**, 2239 (1969)
15. B.K. Jennings, Phys. Lett. **196**, 307 (1987)
16. S. Capstick, A. Švarc, L. Tiator, J. Gegelia, M.M. Giannini, E. Santopinto, C. Hanhart, S. Scherer, T.-S.H. Lee, T. Sato, N. Suzuki, BRAG2007 Workshop SUMMARY. Eur. Phys. J. A **35**(3), 253–266 (2008)
17. N. Kaiser, P.B. Siegel, W. Weise, Phys. Lett. B **362**, 23 (1995)
18. T. Inoue, E. Oset, M.J. Vicente Vacas, Phys. Rev. C **65**, 035204 (2002)
19. M.F.M. Lutz, E.E. Kolomeitsev, Nucl. Phys. A **755**, 29 (2005)
20. M. Döring, Nucl. Phys. A **786**, 164 (2007)
21. O. Krehl, C. Hanhart, S. Krewaldand, J. Speth, Phys. Rev. C **62**, 025207 (2000)
22. A.M. Gasparyan, J. Haidenbauer, C. Hanhart, J. Speth, Phys. Rev. C **68**, 045207 (2003)
23. T.-S.H. Lee, in *Proceedings of the Workshop on the Physics of Excited Nucleons (NSTAR 2005)*, Tallahassee, Florida, 10–15 October 2005 (World Scientific, Singapore, 2006), pp. 1–16
24. N. Mathur, Y. Chen, S.J. Dong, T. Draper, I. Horváth, F.X. Lee, K.F. Liu, J.B. Zhang, Phys. Lett. B **605**, 137 (2005)
25. T. Barnes, F.E. Close, Phys. Lett. B **123**, 221 (1983)
26. Z. Li, Phys. Rev. D **44**, 2841 (1991)
27. C.E. Carlson, N.C. Mukhopadhyay, Phys. Rev. Lett. **67**, 3745 (1991)
28. T. Sato, T.-S.H. Lee, Phys. Rev. C **54**, 2660 (1996)
29. T. Sato, T.-S.H. Lee, Phys. Rev. C **63**, 055201 (2001)
30. L. Tiator, D. Drechsel, S. Kamalov, M.M. Giannini, E. Santopinto, A. Vassallo, Eur. Phys. J. A **19**, 55–60 (2004)
31. A. Matsuyama, T. Sato, T.-S.H. Lee, Phys. Rep. **493**, 193 (2007)
32. I.R. Afnan, Acta Phys. Pol. **8**, 2397 (1998)
33. A.W. Thomas, Adv. Nucl. Phys. **13**, 1 (1984)
34. D.S. Hwang, D.-W. Kim, Phys. Lett. B **601**, 137 (2004)
35. Yu.S. Kalashnikova, Phys. Rev. D **72**, 034010 (2005)
36. P.J. Fink Jr., G. He, R.H. Landau, J.W. Schnik, Phys. Rev. C **41**, 2720 (1990)
37. R.E. Cutkosky, C.P. Forsyth, R.E. Hendrick, R.L. Kelly, Phys. Rev. D **20**, 2839 (1979)
38. M. Batinić, I. Šlaus, A. Švarc, B.M.K. Nefkens, Phys. Rev. C **51**, 2310 (1995)
39. M. Batinić, I. Dadić, I. Šlaus, A. Švarc, B.M.K. Nefkens, T.S.-H. Lee, Phys. Scr. **58**, 15 (1998)
40. W.-M. Yao, et al. (Particle data Group). J. Phys. G: Nucl. Part. Phys. **33**, 1 (2006)
41. S. Ceci, A. Švarc, B. Zauner, D.M. Manley, S. Capstick, Phys. Lett. B **659**, 228 (2008)
42. S. Ceci, J. Stahov, A. Švarc, S. Watson, B. Zauner, Phys. Rev. D **77**, 116007 (2008)
43. B.C. Liu, B.S. Zou, Phys. Rev. Lett. **96**, 042002 (2006)
44. A. Švarc, in *Proceedings of the Workshop on the Physics of Excited Nucleons, NSTAR2005: Plenary Talks*, Tallahassee, 12–15 October 2005 (World Scientific, Singapore, 2006), pp. 37–47

45. http://gwdac.phys.gwu.edu/analysis/pin_analysis.html
46. R.A. Arndt, W.J. Briscoe, I.I. Strakovsky, R.L. Workman, Phys. Rev. C **74**, 045205 (2006)
47. S. Ceci, A. Švarc, B. Zauner, Phys. Rev. Lett. **97**, 062002 (2006)
48. S. Capstick, N. Isgur, Phys. Rev. D **34**, 2809 (1986)
49. B. Saghai, Z. Li, Eur. Phys. J. A **11**, 217 (2001)
50. M.M. Giannini, E. Santopinto, A. Vassallo, Eur. Phys. J. A **12**, 447 (2001)
51. G.-Y. Chen, S. Kamalov, S.N. Yang, D. Drechsel, L. Tiator, Nucl. Phys. A **723**, 447 (2003)
52. W.T. Chiang, S.N. Yang, M. Vanderhaeghen, D. Drechsel, Nucl. Phys. A **723**, 205 (2003)
53. V.A. Tryasuchev, Eur. Phys. J. A **22**, 97 (2004)
54. N.G. Kelkar, M. Nowakowski, K.P. Khemchandani, S.R. Jain, Nucl. Phys. A **730**, 121 (2004)
55. A.V. Anisovich, A. Sarantsev, O. Bartholomy, E. Klempt, V.A. Nikonov, U. Thoma, Eur. Phys. J. A **25**, 427 (2005)
56. A.V. Sarantsev, V.A. Nikonov, A.V. Anisovich, E. Klempt, U. Thoma, Eur. Phys. J. A **25**, 441 (2005)
57. B. Juliá-Díaz, B. Saghai, T.-S.H. Lee, F. Tabakin, Phys. Rev. C **73**, 055204 (2006)