Regge Trajectories of Radial Meson Excitations: Exploring the Dyson-Schwinger–Bethe-Salpeter Approach

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Abstract The combined Dyson-Schwinger and Bethe-Salpeter equations in rainbow-ladder approximation are used to search for Regge trajectories of mesons in the pseudo-scalar and vector channels. We focus on the often employed Alkofer-Watson-Weigel kernel which is known to deliver good results for the ground state meson spectra; it provides linear Regge trajectories in the $J^P = 0^-$ channel.

1 Introduction

Despite of the apparent simplicity of the Lagrangian where Quantum Chromodynamics (QCD) is based upon, it encodes an enormous richness of phenomena, most of them related to the non-perturbative regime. While lattice QCD allows for an access to many facets of the hadron spectra, the so called XYZ states pose still a challenge [1]. Apart the quantitatively adequate description of low-lying hadron states in various flavor channels, the higher excitations call also for a description and confrontation with experimentally well established facts. It is known for a long time that mesons of a given flavor composition can be grouped on Regge trajectories according to $M_n^2 = M_0^2 + n \mu ^2$, where $M_n$ stands for the mass (energy) labeled by the radial quantum number $n = 0, 1, 2, \ldots$, $M_0$ denotes the ground state mass of a respective trajectory and $\mu ^2 = 1.25$ GeV$^2$ [2] or 1.35 GeV$^2$ [3] is a universal slope parameter (cf. [3, 4] for a recent account and [5, 6, 7, 8] for the discussion of the
experimental data base). More generally, Ref. [9] advocates an ordering according to \( M_{nJ} \equiv \hat{a}(n+J) + \hat{c} \), where \( J \) stands for the angular momentum and \( \hat{a} \) and \( \hat{c} \) are appropriate constants, see also [10].

While being a phenomenological ordering scheme, the arrangement of hadron states on Regge trajectories should emerge from QCD, ideally directly without approximations or based on certain symmetries or as result of suitable models. In fact, the relativistic quark model [11, 12] delivers such linear trajectories. Also holographic models based on the AdS/CFT correspondence (cf. [6, 13]) cope with Regge trajectories [14], or even use them as input for constraining the dilaton dynamics for further investigations [15, 16, 17]. Moreover, functional formulations of QCD such as combined Dyson-Schwinger (DS) and Bethe-Salpeter (BS) equations address the issue of recovering Regge trajectories [18, 19] with appropriate interactions kernels and truncation schemes [20]. The latter approach is interesting since it provides the avenue towards addressing the important quest for medium modifications of hadrons in a hot and dense hadron medium [21]. Considering the medium created transiently in the course of relativistic heavy-ion collisions, the interplay of confinement and chiral symmetry restoration poses further challenges [22].

Here, we focus on the question whether the DS-BS approach in rainbow-ladder approximation is capable to deliver Regge trajectories when using simple interaction kernels. To be specific we employ the Alkofer-Watson-Weigel (AWW) kernel [23] in the pseudo-scalar and vector channels and search for the first excited states. Such a study is a prerequisite for the extension to non-zero temperatures [24]. The AWW kernel is known to provide a good description of meson ground states supposed the analytic properties of the quark propagators are properly dealt with [25, 26, 27]. However, in the literature one finds remarks that AWW is less appropriate for a description of excitations due to their sensitivity to long-range interactions [20, 23, 28, 29] (for dedicated studies, cf. [30, 31, 32, 33, 34] for instance). Nevertheless, we feel that a further investigation is timely, in particular w.r.t. the above stressed importance of Regge trajectories as an important feature of the meson spectrum. For the search of meson excitations we employ a novel method of finding solution of the BS equation.

Our paper is organized as follows. In section 2 we recall the DS and BS equations as well as the AWW kernel. Numerical results are described in section 3. We summarize in section 4. The appendix contains some technicalities.

### 2 Recalling the DS and BS equations in rainbow-ladder approximation

The DS equation (also dubbed gap equation) aims at solving

\[
S^{-1}(p) = S_0^{-1}(p) - \int \frac{d^4k}{(2\pi)^4} \left[ -ig^2\gamma^\mu \frac{\tau^a}{2} \right] \Gamma_{\mu\nu}(p,k) \Gamma^{\mu\nu}(p,k) S(k),
\]  

(1)
for the dressed quark propagator $S$, where $S_0$ is the undressed quark propagator, $\gamma^\nu$ are the Dirac matrices with $\{\gamma^\mu, \gamma^\nu\} = 2\delta^\mu\nu$, $\tau^a$ are color matrices, $p$ and $k$ are four-momenta, $g$ is the QCD coupling constant, and $D_{\mu\nu}$ stands for the gluon propagator.

In Euclidean space, the rainbow-ladder approximation means $\Gamma^{\mu\nu}(p,k) \Rightarrow -i\gamma^\mu \frac{e^a}{2}$.

The final equation to be solved reads

$$S^{-1}(p) = S_0^{-1}(p) + \frac{4}{3} \int \frac{d^4k}{(2\pi)^4} [g^2 D_{\mu\nu}(p-k)]^{\gamma^\mu S(k) \gamma^\nu}$$

with $S_0^{-1}(p) = i\not{p} + m_q$, where the parameter $m_q$ is flavor dependent.

The dressed quark propagator $S(p)$ enters the BS equation for the vertex function

$$\Gamma(P,p,k) = \int \frac{d^4k}{(2\pi)^4} K(P,p,k)S(k+\eta_1 P)\Gamma(P,k)S(k-\eta_2 P),$$

where $P$ and $p$ are the total and relative momenta of quarks and $\eta_1 + \eta_2 = 1$ describes momentum sharing ($P \sim (M_{12}, 0)$ for a meson at rest), and the rainbow-ladder approximation for the kernel function

$$K(P,p,k) = -g^2 D_{\mu\nu}(p-k)\left(\gamma^\mu \frac{e^a}{2}\right)\left(\gamma^\nu \frac{e^a}{2}\right).$$

In the Euclidean space, the used BS equation becomes then

$$\Gamma(P,p) = -\frac{4}{3} \int \frac{d^4k}{(2\pi)^4} \gamma^\mu S(k+\eta_1 P)\Gamma(P,k)S(k-\eta_2 P)\gamma^\nu [g^2 D_{\mu\nu}(p-k)].$$

We employ here the AWW kernel, i.e. $D(k^2) \Rightarrow D^{\text{AWW}}(k^2)$ in the decomposition of the gluon propagator in Landau gauge, $g^2 D_{\mu\nu}(k) = (\delta_{\mu\nu} - k_{\mu}k_{\nu}k^2)D(k^2)$ with

$$D^{\text{AWW}}(k^2) = \frac{4\pi^2 Dk^2}{\omega^2} e^{-\frac{k^2}{\omega^2}},$$

with the interaction strength parameter $D$ and the interaction range parameter $\omega$. (In what follows, we employ $\eta_1 = \eta_2 = \frac{1}{2}$ and the standard model parameters $\omega = 0.5\text{GeV}$ and $D = 16\text{GeV}^{-2}$, unless explicitly noted.) It is the IR part of the Maris-Tandy kernel [35].

3 Numerical methods and results

The numerical details for solving the above quoted DS and BS equations with given truncations and approximations are described in [36]. In vacuum, the quark propagator can be decomposed as $S^{-1}(p) = i\not{p} A(p) + B(p)$ to split (2) into two coupled integral equations for $A$ and $B$ which are needed for complex arguments $p$. Figure[1]
exhibits examples for these functions for positive, real values of the momentum $p$. Note the non-linear dependence on the quark mass parameter $m_q$.

![Graphs of A(p) and B(p)](image)

Fig. 1: Propagator functions $A(p)$ (left) and $B(p)$ (right) on the positive real axis. Red curve: $m_q = 0$ MeV (chiral limit), blue curve: $m_q = 5$ MeV, yellow curve: $m_q = 115$ MeV, green curve: $m_q = 1130$ MeV.

Making an expansion of the BS vertex function \( S \) into spin-angular functions, spherical harmonics and Gegenbauer polynomials one arrives at a matrix equation

\[
X_\alpha = S_{\alpha\beta} X_\beta
\]

with $\alpha, \beta = 1, \cdots N$, where $N = \alpha_{\text{max}} N_{\text{Gegenbauer}} N_{\text{Gauss}}$. Here, $\alpha_{\text{max}}$ denotes the number of spin-angular harmonics, $N_{\text{Gegenbauer}}$ is the number of included Gegenbauer polynomials, and $N_{\text{Gauss}}$ stands for the mesh number of a Gaussian integration over internal momenta. The chain of manipulations that leads to the quantity $S_{\alpha\beta}$ is recalled in the Appendix, where also the elements of $X$ are defined.

The energy of mesons as $\bar{q}q$ bound states is determined by $\det(S - I) = 0$ with $S$ being a function of $M$. An example is exhibited in Fig. 2.

![Smooth determinant function](image)

Fig. 2: Smooth determinant function $\det(S - I)$ as a function of $M_{12}$ for the pion channel ($m_{1,2} = m_q = m_u = 5$ MeV). For $\omega = 0.3$ GeV and $D = 205.761$ GeV$^{-2}$. The arrows denote the masses of ground state (g.s.), first excited state (1st), second excited state (2nd), third and fourth excited state (3rd, 4th).

The AWW kernel depends on two parameters, $D$ and $\omega$; in addition the quark masses $m_{1,2}$ ($m_1 = m_2 = m_q$ for equal quark-mass mesons) must be adjusted. Figure 3 exhibits examples for the pseudo-scalar (left) and vector (right) channels.

Freezing in $m_u = m_d = 5$ MeV, $m_s = 115$ MeV, $m_c = 1130$ MeV and keeping our standard parameters $\{\omega, D\}$ we find in the $J^p = 0^-$ channel (in units of GeV, ex-
experimental values in parentheses) \( M_\pi = 0.137 (0.140), 0.986 (1.300), 1.369 (1.812), \)
\( M_K = 0.492 (0.494), 1.162 (1.460), M_{\bar{s}s} = 0.693 (-), 1.278 (-), 1.572 (-), \)
\( M_D = - (1.870), M_{D_s} = 2.075 (1.968), 2.313 (-), M_{\eta_s} = 2.984 (2.984), 3.278 (3.639), 3.557 (-), \)
where a “-” means either no state found or missing information. The results for the \( J^{PC} = 1^{--} \) states are \( M_\rho = 0.758 (0.775), 1.041 (1.465), 1.287 (1.720), \)
\( M_K^* = 0.945 (894), 1.264 (1.414), M_\phi = 1.077 (1.019), 1.402 (1.680), 1.598 (2.175), \)
\( M_{D_s} = - (2.010), M_{D_s} = - (2.112), M_{J/\psi} = 3.136 (3.097), 3.346 (3.686), 3.593 (3.773). \)
In addition, we note \( f_\pi = 0.133 \text{ GeV} \) and \( \langle \bar{q}q \rangle = (-0.255 \text{ GeV})^3 \). The overwhelming impression is that, despite of the truncation and the simple AWW kernel, quite reasonable numbers are delivered, however, with some drastic deviations, e.g.
the pure pseudo-scalar \( \bar{s}s \) states do not appear in nature.

Another issue is the discrepancy of calculated and experimental values of excitations, e.g. in the pion channel. Instead of disputing the impact of mixing effects, we take our calculated values and check the arising Regge trajectories for linearity, see Fig. 4 for an example with apparently linear trajectory. Some survey is exhibited in Fig. 5 where a few 0− states are depicted (left column) and the Regge slope parameters as well as a linearity measure are displayed too (right column). In that study, we freeze in \( \alpha = \omega \bar{s}s \) and vary the parameter \( \omega \). As known, the ground state masses are kept constant under such a variation, but evidently the excited states depend on \( \omega \), even up to a disappearance of certain states, e.g. \( \pi \) and \( \bar{s}s \). The slope changes

Fig. 3: Contour plot of pseudo-scalar (left) and vector (right) meson ground state masses in units of GeV for varying quark masses \( m_1 \) and \( m_2 \). The colored bullets denote the experimental values of meson ground states (red: \( \pi/\rho \), green: \( K/\phi \), violet: \( \eta_{c}/J_{\psi} \)) which could be used for extracting the bare quark mass parameters \( m_{1,2} \) (vertical and horizontal dashed lines, labeled by the corresponding quark flavor); the corresponding value of \( m_c = 1.110 \text{ GeV} \) in the vector channel can be compared with one suggested in the pseudo-scalar channel, with optimum value \( 1.130 \text{ GeV} \).

In the white region, no solutions of the BS equation could found w/o accounting explicitly for the pole structure of \( \bar{s}s \) in the complex momentum plane.
with $\omega$, while the linearity is strikingly good. This is in contrast to the $1^{-+}$ channel (not displayed; for details cf. \cite{35}), where, at fixed values of $\omega^3D$, also the ground states vary with changing $\omega$; linear Regge trajectories are hardly accessible within the preferred parameter range adjusted to $0^-$ states.

4 Summary

In summary we test the capability to catch the first excited states of mesons (pseudo-scalar and vector channels) by using the Dyson-Schwinger and Bethe-Salpeter equations in rainbow-ladder approximation equipped with the Alkofer-Watson-Weigel kernel. This is the first step of an attempt to describe hadron properties, and thus implicitly confinement and relevant scales, together with the extension to finite temperatures and baryon densities in follow-up investigations. According to our contemporary understanding, at some temperature and at small baryon density, hadrons as quasi-particle degrees of QCD should disappear in favor of quasi-quark and quasi-gluon degrees of freedom. The transition happens gradually and may depend on the flavor channel under consideration. For larger baryon densities, the transition could be abrupt, supposed a critical point occurs in the phase diagram of strongly interacting matter. A few large-scale heavy-ion experiments, e.g. the beam-energy scan at RHIC, NA61/SHINE at SPS, CBM at SIS100, as well at NICA and J-PARC address in their physics programs the critical point search. For that, both the properties of hadrons as individual entities and the behavior of hadron matter are key quantities in reconstructing the final state of strong-interaction matter in collision.

Acknowledgements The authors gratefully acknowledge the collaboration with S. M. Dorkin, T. Hilger and M. Viebach on the topic.

Appendix: Spin-angular harmonics

The BS vertex function $\Gamma$ can be expanded into spin-angular harmonics:
Fig. 5: Dependence of the spectrum of pions (a: $m_q = m_u$), kaons (b: $m_1 = m_u, m_2 = m_s$), fictitious pseudo-scalar $\bar{s}s$ states (c: $m_q = m_s$), and $\eta_c$ (d: $m_q = m_c$) as a function of $\omega$ for $a = 0.5\text{GeV}^3$ (left column). The right column depicts the corresponding Regge slope coefficients $\beta$ (empty violet squares), the quadratic term $c$ (empty green squares) in fits of the spectra by $M_n^2 = M_0 + \beta n + cn^2$, $\beta \equiv \mu^2$, and the deviation measure from linear behavior $|c/\beta|$ (filled blue squares).
\[ \Gamma(p) = \sum_{\alpha=1}^{\alpha_{\text{max}}} \Gamma_\alpha(p) = \sum_{\alpha=1}^{\alpha_{\text{max}}} g_\alpha(p) \mathcal{F}_\alpha(p) \]  

with functions \( g_\alpha \) fulfilling the orthogonality relation \( g_\alpha(p) = \int d\Omega_p \text{Tr}[\Gamma(p) \mathcal{F}^\dagger_\alpha(p)] \).

For pseudo-scalar mesons (\( J^{PC} = 0^{-+} \)), the number of independent spin-angular harmonics \( \alpha_{\text{max}} = 4 \), and the set is chosen as

\[
\begin{align*}
\mathcal{F}_1(p) &= \frac{1}{\sqrt{16\pi}} \gamma^5 = \mathcal{F}^\dagger_1(p), \\
\mathcal{F}_2(p) &= \frac{1}{\sqrt{16\pi}} \gamma^\rho \gamma^5 = -\mathcal{F}^\dagger_2(p), \\
\mathcal{F}_3(p) &= -\frac{1}{\sqrt{16\pi}} n_p \gamma^\rho \gamma^5 = \mathcal{F}_3(p), \\
\mathcal{F}_4(p) &= -\frac{1}{\sqrt{16\pi}} n_p \gamma^\rho = \mathcal{F}_4(p),
\end{align*}
\]

and for vector mesons (\( J^{PC} = 1^{+-} \)), \( \alpha_{\text{max}} = 8 \) with

\[
\begin{align*}
\mathcal{F}_1(p) &= \sqrt{\frac{1}{16\pi}} \xi_{\mu\nu} = \mathcal{F}^\dagger_1(p), \\
\mathcal{F}_2(p) &= -\sqrt{\frac{1}{16\pi}} \gamma^\rho \xi_{\mu\nu} = \mathcal{F}^\dagger_2(p), \\
\mathcal{F}_3(p) &= -\sqrt{\frac{3}{16\pi}} (n_p \xi_{\mu\nu}) = \mathcal{F}_3(p), \\
\mathcal{F}_4(p) &= \sqrt{\frac{3}{32\pi}} \gamma^\rho [-(n_p \xi_{\mu\nu}) + \#_p \xi_{\mu\nu}] = -\mathcal{F}_4(p), \\
\mathcal{F}_5(p) &= \sqrt{\frac{1}{32\pi}} \#_p \xi_{\mu\nu} + 3(n_p \xi_{\mu\nu}) \#_p = -\mathcal{F}_5(p), \\
\mathcal{F}_6(p) &= \sqrt{\frac{3}{32\pi}} \gamma^\rho [\xi_{\mu\nu} + 3(n_p \xi_{\mu\nu}) \#_p] = \mathcal{F}_6(p), \\
\mathcal{F}_7(p) &= -\sqrt{\frac{3}{16\pi}} \gamma^\rho (n_p \xi_{\mu\nu}) = \mathcal{F}_7(p), \\
\mathcal{F}_8(p) &= \sqrt{\frac{3}{32\pi}} [-(n_p \xi_{\mu\nu}) + \#_p \xi_{\mu\nu}] = -\mathcal{F}_8(p).
\end{align*}
\]

Scalar products are displayed here in Minkowski space; \( n_p \) is the unit vector \( n_p = (0, p/|p|) \), \( \xi_{\mu\nu} = (0, \xi_{\mu\nu}) \) is the polarization vector fixed by \( \xi_{+1} = -(1, i, 0)/\sqrt{2}, \) \( \xi_{-1} = (1, -i, 0)/\sqrt{2}, \xi_0 = (0, 0, 1) \) and slashed quantities such as \( \mathcal{F}^\dagger \) represent \( \gamma^\mu x_\mu \).

The partial amplitudes \( \Gamma_\alpha(p) \) and the interaction kernel (6) are decomposed over the basis of spherical harmonics \( Y_{lm}(\theta, \phi) \) and normalized Gegenbauer polynomials \( X_{al}(\chi) \).

\[ \mathcal{Z}_{al;m} = X_{al}(\chi)Y_{lm}(\theta, \phi) \]

\[ = \sqrt{\frac{2l+1}{\pi}} \frac{(n+1)(n-1)!!}{(n+l+1)!} \sin^{l+1} \chi C_{n-l}^{l+1} \cos \chi Y_{lm}(\theta, \phi), \]

with familiar Gegenbauer polynomials \( C_{n-l}^{l+1} \cos \chi \). The hyper angle \( \chi \) is defined by \( \cos \chi = p_4/\tilde{p} \) and \( \sin \chi = |p|/\tilde{p}, \) where \( \tilde{p} = (p_4^2 + p_5^2)^{1/2} \) is the modulus for
an Euclidean four-vector \( p = (p_4, \mathbf{p}) \). The partial decompositions of \( \Gamma_\alpha(p) \) and \( D^{\text{AWW}}(p-k) \) read

\[
\Gamma_\alpha(p) = \sum_n \varphi^n_{\alpha,l_\alpha}(\mathbf{p}) X_{n\alpha}(\chi_\mathbf{p}) \mathcal{T}_\alpha(\mathbf{p}), \tag{12}
\]

\[
D^{\text{AWW}}(p-k) = 2\pi^2 \sum_{k,\lambda} \frac{1}{\kappa+1} V_k(\mathbf{p}, \mathbf{k}) X_{k\lambda}(\chi_\mathbf{p}) X_{k\lambda}(\chi_\mathbf{k}) Y_{\lambda\mu}(\Omega_\mathbf{p}) Y^{*}_{\lambda\mu}(\Omega_\mathbf{k}), \tag{13}
\]

where \( V_k(\mathbf{p}, \mathbf{k}) \) are the partial kernels and \( \varphi^n_{\alpha,l_\alpha}(\mathbf{p}) \) are the expansion coefficients of the partial amplitudes. Actually \( l_\alpha \) is restricted by the corresponding orbital momentum encoded in \( \mathcal{T}_\alpha(\mathbf{p}) \). For \( \mathcal{T}_{1,2}(\mathbf{p}) \) from Eq. \( (8) \), \( l_\alpha = 0 \) holds, while for \( \mathcal{T}_{3,4}(\mathbf{p}) \) \( l_\alpha = 1 \). In analogy for vector mesons (see \( (9) \)), \( l_\alpha = 0 \) for \( \mathcal{T}_{1,2}(\mathbf{p}) \), \( l_\alpha = 1 \) for \( \mathcal{T}_{3,4,7,8}(\mathbf{p}) \) and \( l_\alpha = 2 \) for \( \mathcal{T}_{5,6}(\mathbf{p}) \).

Changing the integration variables to the hyperspace, \( d^4k = \bar{k}^3 \sin^2 \chi_k \sin \theta_k d\bar{k} \times d\chi_k d\theta_k d\phi_k \), inserting Eq. \( (12) \) and \( (13) \) into \( (5) \) and performing the necessary angular integration, a system of integral equations for the expansion coefficients \( \varphi^n_{\alpha,l_\alpha}(\mathbf{p}) \) as the \( N \) elements of \( X \) remains:

\[
\varphi^n_{\alpha,l_\alpha}(\mathbf{p}) = \sum_{n}^{\infty} \sum_{\beta} \int d\bar{k}^3 S_{\alpha\beta}(\mathbf{p}, \mathbf{k}, \mathbf{n}) \varphi^n_{\beta,l_\beta}(\mathbf{k}). \tag{14}
\]

The explicit expression for \( S_{\alpha\beta}(\mathbf{p}, \mathbf{k}, \mathbf{n}) \) reads

\[
S_{\alpha\beta}(\mathbf{p}, \mathbf{k}, \mathbf{n}) = \sum_{\kappa} \int \sin^2 \chi_k d\chi_k X_{k\alpha}(\chi_\mathbf{k}) X_{k\beta}(\chi_\mathbf{k}) \sigma_{\chi_\mathbf{n}}(\bar{k}^2) \sigma_{\chi_\mathbf{n}}(\bar{k}^2) \times A_{\alpha\beta}(\mathbf{p}, \mathbf{k}, \mathbf{n}), \tag{15}
\]

where \( \bar{k}^2 \) is given by \( \bar{k}^2 = k^2 - \eta^2 \eta_{1,2} M^2_{1,2} \pm \eta_{1,2} M_{1,2} \cos \chi \) with momentum partitioning parameters \( \eta_1 + \eta_2 = 1 \) and \( A_{\alpha\beta}(\mathbf{p}, \mathbf{k}, \mathbf{\chi_n}) \) results from calculations of traces and angular integrations as

\[
A_{\alpha\beta}(\mathbf{p}, \mathbf{k}, \mathbf{\chi_n}) = \int \sin^2 \chi_k d\chi_k d\Omega_k d\Omega_k V_{\alpha}(\mathbf{p}, \mathbf{k}) X_{\alpha\beta}(\chi_\mathbf{p}) X_{\alpha\beta}(\chi_\mathbf{k}) Y_{\lambda\mu}(\Omega_\mathbf{p}) Y_{\lambda\mu}^{*}(\Omega_\mathbf{k}) \times \text{Tr}[d\mu^\alpha((p-k)^2) \gamma^\mu \cdots \mathcal{T}_\alpha(\mathbf{p}) \cdots \mathcal{T}_\alpha(\mathbf{p}) \gamma^\nu]. \tag{16}
\]

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