Charmed baryons in bootstrap quark model

S.M. Gerasyuta*, D.V. Ivanov

Department of Theoretical Physics,
St. Petersburg State University
198904, St.Petersburg, Russia

Abstract

In the framework of dispersion relation technique the relativistic three-quark equations including heavy quarks are found. The approximate solutions of the relativistic three-particles equations based on the extraction of leading singularities of amplitudes are obtained. The mass values of S-wave multiplets of charmed baryons are calculated.

1 Introduction

After detection of the $J/\psi$-particles, the study of heavy quark properties has become one of leading directions in particle physics. Some features of $J/\psi$-meson and other states of charmonium stimulated the development of theory of strong interactions (QCD). The further progress is connected with detection of the bottomonium (family Y) and B-mesons (particles containing fifth b-quark). At present number of experimental and theoretical work is carried out, on the basis of which the uniform conception about dynamics of bound states of heavy quarks were obtained. The most reasonable and adequate approaches are the method of dispersion sum rules in QCD [1-5] and non-relativistic potential model [6-12], in the framework of which the quantitative description practically of all known properties of quarkonium ($Q\bar{Q}$) was received.

Doubtless advantages of potential approach to description of heavy mesons are their simplicity and presentation. It permits to calculate (in the framework of accepted model of potential) position of levels, width of radiation transitions between levels and width of annihilation decays. Thus, one can find actually the behaviour of system in enough wide interval of distances, that is important for understanding of dynamics of quark interactions. The serious difficulty of potential approach is that it is impossible completely to be justified in the framework of QCD.

*Present address: Department of Physics, LTA, Institutski Per.5, St.Petersburg 194021, Russia
The method of sum rules of QCD [1,2], based on complex non-perturbative structure of QCD vacuum was successfully used for the study of charmonium. To realize the spectrum of charmonium, it is necessary to know the quark masses and the nature of the connecting forces. According to the method of sum rules of QCD, quarks exist in the complex medium – non-perturbative vacuum, which densely populated by long-wave fluctuations of gluon field. These non-perturbative fluctuations determine the effective quark attraction. The sum rules of QCD permit to calculate the effect of long-wave gluon fluctuations and to make exact predictions for lowest levels with a different quantum numbers in terms of fundamental parameters. The most considerable difficulty of application of the sum rules method to the charmonium is account of the relativistic corrections. It is necessary to take into account with purely relativistic effects (described by Breit-Fermi hamiltonian) radiative corrections of higher orders. Such calculation is so far away.

The problem of the charmed baryons is investigated considerably less. The sum rules of QCD were applied for the first time to the charmed baryons in the paper [13]. However consideration was carried out within the limits of infinite mass of c-quark and also the contribution of the gluon condensate was not taken into account in this case. The subsequent calculations have allowed to calculate the masses of the S-waves charmed baryons with one c-quark [14,15]. The potential quark models enable to calculate the mass spectrum of the charmed baryons of multiplets $J^P = \frac{1}{2}^+, \frac{3}{2}^+$ [16-26]. The most considerable difficulty of application of these methods is account of relativistic corrections at consideration of system of heavy and light quarks.

The present paper is devoted to calculation of mass spectrum of charmed baryons in the framework of the bootstrap quark model. The relativistic Faddeev equations are constructed in the form of the dispersion relations over the two-body subenergy. The approximate solution of integral three-quark equations are obtained by taking into account two-particle and triangle singularities, all the weaker ones being neglected. The mass values of charmed low-lying multiplets $J^P = \frac{1}{2}^+, \frac{3}{2}^+$ are calculated.

In section 2 the relativistic three-particle equations, which describe the interaction of light and heavy quarks in the charmed baryons are obtained. Section 3 is devoted to discussion of calculation results of the low-lying charmed baryons mass spectrum. In Conclusion we discuss the status of the considered model.

In Appendix A the three-particle integral equations for the charmed baryons of multiplets $J^P = \frac{1}{2}^+, \frac{3}{2}^+$ are given. In Appendix B we present the approximate algebraic equations for the charmed baryons, accounting only the contribution of two-particle and triangle singularities in the scattering amplitude.
Three-particle quark amplitudes of the S-wave charmed baryons

In the papers [27,28] relativistic integral three-particle Faddeev equations was constructed in the form of the dispersion relations over the two-body subenergy. Using the method of extraction of the main singularities of amplitude the mass spectrum of S-wave baryons, including u, d, s-quarks is calculated. In present paper this method is applied for study of the mass spectrum of baryons, including c-quarks.

For the clarity we derive the relativistic Faddeev equation for the particle $\Sigma^+(J_P = \frac{1}{2}^+)$, $c\Sigma^+(J_P = 0^+ , \frac{1}{2}^+)$, $(J_P = 0^+ , 1^+)$, in Fig.1 a. Subsequent two-quark interactions lead to the diagrams, shown in Fig.1 b-f. These diagrams might be divided into three groups so in each group the same pair of particles undergoes the last interaction (Fig.2). Consideration of diagram in Fig.1 allows to obtain the equation for three-particle amplitudes $A_J(s, s_{ik})$, $(J_P = 0^+, 1^+)$ in the graphic form (Fig.3). Here the invariants $s_{ik} = (p_i + p_k)^2$, $s = (p_i + p_j + p_k)^2$ for the quarks with momenta $p_i, p_j, p_k$ are given.

In order to present the amplitude $A_J(s, s_{ik})$ in the form of the dispersion relation explicitly, it is necessary to define the amplitude of two-quark interaction $a_J(s_{ik})$. We use the results of bootstrap quark model [29] and write down the diquark amplitude $a_J(s_{ik})$ in the following way:

$$a_J(s_{ik}) = \frac{G_J^2(s_{ik})}{1 - B_J(s_{ik})}, \quad (1)$$

$$B_J(s_{ik}) = \int_{(m_i + m_k)^2}^{\infty} \frac{ds_{ik}' \rho_J(s_{ik}') G_J^2(s_{ik}')} {s_{ik}' - s_{ik}} \quad (2)$$

$$\rho_J(s_{ik}) = \frac{(m_i + m_k)^2}{4\pi} \left[ \gamma(J_P) \left( \frac{s_{ik}}{(m_i + m_k)^2} \right)^2 + \beta(J_P) + \delta(J_P) \right] \times \sqrt{\frac{s_{ik} - (m_i + m_k)^2}{s_{ik}}} \quad (3)$$

Here $G_J$ is the vertex function of diquark, $B_J(s_{ik})$ is the Chew-Mandelstam function [30], $\rho_J(s_{ik})$ is the phase space of diquark (see Table 1).

In the case in question the interacting quarks do not produce the bound state, then the integration in the dispersion integrals (4)-(6) is carried out from $(m_1 + m_2)^2$ to $\infty$. To write
Here $\alpha$ state and that of the particle 3 in the final state, which is taken in the c.m. of the particles diquark.

The index "c" shows the presence of one (c) or two (cc) charmed quarks in the system of equations (4)-(6) (Fig.3), it is necessary to account the spin-flavour part of the charmed baryon $\Sigma^+_c$ wave function: $|\Sigma^+_c\rangle = \sqrt{\frac{3}{5}}\{u \uparrow d \uparrow c \downarrow\} - \frac{1}{\sqrt{5}}\{u \uparrow d \downarrow c \uparrow\} - \frac{1}{\sqrt{5}}\{u \downarrow d \uparrow c \uparrow\}$. As a result we receive the system of integral equations:

$$A_1(s, s_{12}) = \hat{a} b_1(s_{12})L_1(s_{12}) + K_1(s_{12}) \left[ \frac{1}{4} A_1^{(c)}(s, s'_{13}) + \frac{3}{4} A_0^{(c)}(s, s'_{13}) + \frac{1}{4} A_1^{(c)}(s, s'_{23}) + \frac{3}{4} A_0^{(c)}(s, s'_{23}) \right]$$

(4)

$$A_1^{(c)}(s, s_{13}) = \hat{a} b_1^{(c)}(s_{13})L_1^{(c)}(s_{13}) + K_1^{(c)}(s_{13}) \left[ \frac{1}{2} A_1(s, s'_{12}) + \frac{3}{4} A_0(s, s'_{12}) - \frac{1}{4} A_1(s, s'_{12}) + \frac{1}{2} A_1(s, s'_{23}) + \frac{3}{4} A_0(s, s'_{23}) - \frac{1}{4} A_1(s, s'_{23}) \right]$$

(5)

$$A_0^{(c)}(s, s_{23}) = \hat{a} b_0^{(c)}(s_{23})L_0^{(c)}(s_{23}) + K_0^{(c)}(s_{23}) \left[ \frac{1}{2} A_1(s, s'_{13}) + \frac{1}{4} A_1^{(c)}(s, s'_1) + \frac{1}{4} A_0^{(c)}(s, s'_{13}) + \frac{1}{2} A_1(s, s'_{12}) + \frac{1}{4} A_1^{(c)}(s, s'_1) + \frac{1}{4} A_0^{(c)}(s, s'_{12}) \right]$$

(6)

In eq.(4)-(6) we introduce functions

$$L_J(s_{ik}) = \frac{G_J(s_{ik})}{1 - B_J(s_{ik})}$$

(7)

and integral operators

$$K_J(s_{ik}) = L_J(s_{ik}) \int_{m_1}^{\infty} \frac{ds'_{ik} \rho_J(s'_{ik}) G_J(s'_{ik})}{s'_{ik} - s_{ik}} \int_{-\infty}^{\infty} \frac{dz}{2}$$

(8)

Here $b_J(s_{ik}) = \int_{m_1}^{\infty} \frac{ds'_{ik} \rho_J(s'_{ik}) G_J(s'_{ik})}{s'_{ik} - s_{ik}}$ is truncated Chew-Mandelstam function, $z$ is cosine of the angle between the relative momentum of the particles 1 and 2 in the intermediate state and that of the particle 3 in the final state, which is taken in the c.m. of the particles 1 and 2. The coefficient $\hat{a}$ corresponds to the vertex of three-quark creation (later its value is not used). The index "c" shows the presence of one (c) or two (cc) charmed quarks in diquark.

Invariants $s'_{13}$ and $s'_{12}$ for various quarks are connected by the relation:

$$s'_{13} = m_1^2 + m_3^2 - \frac{(s'_{12} + m_3^2 - s)(s'_{12} + m_1^2 - m_2^2)}{2s'_{12}} \pm$$

$$\pm \frac{z}{2s'_{12}} [(s'_{12} - (m_1 + m_2)^2)(s'_{12} - (m_1 - m_2)^2)]^{1/2} \times$$

(9)
on the extraction of the leading singularities which are close to the region singularities. The main singularities in Amplitudes with different number of rescattering (Fig.1) have the following structure of $\alpha$ for reduced amplitudes essentially the calculated mass spectrum. Similarly two other equations of system (5), (6) also the "soft" cutting, for instance $G$ off parameter $\Lambda J$ to replace $\sqrt{A}$ are given in Appendix A.  

We extract the diquark singularity of the amplitude $A_J(s, s_{ik})$:

$$A_J(s, s_{ik}) = \frac{\alpha_J(s, s_{ik}) b_J(s_{ik}) G_J(s_{ik})}{1 - B_J(s_{ik})}$$  \hspace{3cm} (10)$$

Then, for example, first equation of system (4)-(6) is possible to be rewritten for the amplitude $\alpha_1(s, s_{12})$ in the form:

$$\alpha_1(s, s_{12}) = \hat{a} + \frac{1}{b_1(s_{12})} \int_{(m_1 + m_2)^2}^{\Lambda_1(1,2)} \frac{ds_{12}'}{\pi} \frac{\rho_1(s_{12}') G_1(s_{12}')} {s_{12}' - s_{12}} \times$$

$$\times \int_{-1}^{1} dz \frac{G_1^{(c)}(s_{13}') b_1^{(c)}(s_{13}') \frac{1}{2} \alpha_1^{(c)}(s, s_{13}') + G_0^{(c)}(s_{13}') B_0^{(c)}(s_{13}') \frac{3}{2} \alpha_0^{(c)}(s, s_{13}')} {1 - B_0^{(c)}(s_{13}')}}$$  \hspace{3cm} (11)$$

The formula for $s_{23}'$ is similar to (7) with $z$ replaced by $-z$. Therefore in (4) it is possible to replace $[A_1^{(c)}(s, s_{13}') + A_1^{(c)}(s, s_{23}')]$ by $2A_1^{(c)}(s, s_{13}')$. In equation (11) we introduce the cut-off parameter $A_J(i, k)$ at large value of $s_{12}$. We choose the "hard" cutting, but we can use also the "soft" cutting, for instance $G_1(s_{12}) = G_1 \exp(-(s_{12} - 4m^2)^2/\Lambda_1^2)$, and do not change essentially the calculated mass spectrum. Similarly two other equations of system (5), (6) for reduced amplitudes $\alpha_1^{(c)}(s, s_{13})$, $\alpha_0^{(c)}(s, s_{23})$ are obtained (Appendix B).

The construction of the approximate solution of the system of equations (4) - (6) is based on the extraction of the leading singularities which are close to the region $s_{ik} = (m_i + m_k)^2$. Amplitudes with different number of rescattering (Fig.1) have the following structure of singularities. The main singularities in $s_{ik}$ arise from pair rescattering of the particles $i$ and $k$. First of all there are threshold square root singularities. Also possible a pole singularities which correspond to the bound states. Figs.1b,c have only such two-particle singularities. The diagrams in Fig.1d,e apart from two-particle singularities have their own specific triangle singularities. The diagram in Fig.1f describes still more three-particle singularities. Apart from the two-particle and triangle singularities it has its own singularity. It is weaker than singularities of the diagrams Fig.1b-e.

Such classification allows us to search the approximate solution of equations (4) - (6) by taking into account some definite number of leading singularities and neglecting all the weaker ones. We consider the approximation, which corresponds to the single interaction of all three particles (two-particle and triangle singularities).
The contours of integration in complex plane of variable $s'_{13}$ are given in Fig.4. At $(m_1 + m_2)^2 \leq s'_{12} \leq (\sqrt{s} + m_3)^2$ the integration performed over the contour 1, and at $(\sqrt{s} + m_3)^2 \leq s'_{12} \leq \Lambda_J(1,2)$ over the contour 2.

The function $\alpha_J(s,s_{12})$ is a smooth function of $s_{12}$ as compared with the singular part of the amplitude, hence it can be expanded in a series in the singularity point and only the first term of this series should be employed further. It is convenient to take the middle point of Dalitz-plot region which corresponds to $z = 0$. Then one determines the function $\alpha_J(s,s_{12})$ and truncated Chew-Mandelstam function $b_J(s_{12})$ at the point $s_{12} = s_0 = (s + m_1^2 + m_2^2 + m_3^2)/(m_1^2 + m_2^2 + m_3^2)$, where $m_{ik} = \frac{1}{2}(m_i + m_k), i = 1,2,3$. Such a choice of point $s_0$ allows us to replace the system of integral equations (4)-(6) by the algebraic system of equations for $\Sigma^+_c$:

$$\alpha_1(s,s_0) = \hat{\alpha} + \frac{1}{2} I_{1,1}(s,s_0) \frac{b_1^{(c)}(s_0)}{b_1^{(c)}(s_0)} \alpha_1^{(c)}(s,s_0) + \frac{3}{2} I_{1,0}(s,s_0) \frac{b_1^{(c)}(s_0)}{b_0^{(c)}(s_0)} \alpha_0^{(c)}(s,s_0) + \frac{1}{2} I_{0,1}(s,s_0) \frac{b_1^{(c)}(s_0)}{b_0^{(c)}(s_0)} \alpha_1^{(c)}(s,s_0)$$

$$- \frac{1}{2} I_{1,1}(s,s_0) \alpha_1^{(c)}(s,s_0) + \frac{3}{2} I_{1,0}(s,s_0) \frac{b_1^{(c)}(s_0)}{b_0^{(c)}(s_0)} \alpha_0^{(c)}(s,s_0)$$

$$\alpha_0^{(c)}(s,s_0) = \hat{\alpha} + \frac{1}{2} I_{0,1}(s,s_0) \frac{b_1^{(c)}(s_0)}{b_0^{(c)}(s_0)} \alpha_1^{(c)}(s,s_0) + \frac{1}{2} I_{0,0}(s,s_0) \alpha_0^{(c)}(s,s_0)$$

Here

$$I_{J_1, J_2}(s, s_0) = \int_{(m_i + m_k)^2}^{\Lambda_{J_1}(i,k)} \frac{d s'_{ik}}{\pi} \frac{\rho_{J_1}(s'_{ik})}{s'_{ik} - s_0} G_{J_1} \int_{-1}^{1} \frac{dz}{2} \frac{G_{J_2}}{1 - B_{J_2}(s'_{ij})}$$

In our approximation vertex functions $G_{J_2}(s'_{ik})$ are the constants.

The function $I_{J_1, J_2}$ takes into account singularities which corresponds to the simultaneous vanishing of all propagators in the triangle diagrams like those in Fig.1d,e. The solution of the system of equations (12) - (14) permits to determine the pole in $s$ (zero of determinant) which corresponds to the bound state of the three quarks.

Similarly to the case of $\Sigma^+_c$–baryon it is possible to obtain the rescattering amplitudes for all S-wave charmed baryons with $J^P = \frac{1}{2}^+, \frac{3}{2}^+$ including $u$, $d$, $s$, $c$-quarks, which satisfy the systems of integral equations (Appendix A). If we consider the approximation, which corresponds to taking into account two-particle and triangle singularities, and define all the functions of the subenergy variables in the middle point of the physical region of Dalitz-plot $s_0$ then the problem reduces to one of solving simple algebraic system equations (Appendix B).
In the calculation in question the quark masses \( m_u = m_d = m, m_s, m_c \) are not fixed. We fix such quark mass values, which allows us to calculate the masses of charmed baryons using the vertex constants and the cut-off parameters:

\[
g^{(ik)}_J = \frac{m_i + m_k}{2\pi} G^{(ik)}_J, \quad \lambda^{(ik)}_J = \frac{4\Lambda_J(i,k)}{(m_i + m_k)^2}, \quad s_0 = \frac{4s_{ik}}{(m_i + m_k)^2},
\]

where \( J = 0, 1 \); \( m_i \) and \( m_k \) are quark masses in the intermediate state of the corresponding quark loop.

### 3 Calculation results

The masses of S-wave charmed baryons \( J^P = \frac{1}{2}^+, \frac{3}{2}^+ \) including light u, d, s and heavy c-quarks are obtained. Two dimensionless cut-off parameters \( \lambda_q = 10.7 \) and \( \lambda_c = 6.5 \) are chosen for light quarks and c-quark respectively. For the diquarks, which include one light and one heavy quark, the cut-off parameter is equal \( \lambda_{qc} = \frac{1}{4}(\sqrt{\lambda_c} + \sqrt{\lambda_q})^2 \). For light quarks the calculated values of vertex constants \( g_1 = 0.56 \) and \( g_0 = 0.70 \) have coincided with ones, obtained earlier in the bootstrap quark model for light mesons and baryons [27-29]. This coincidence is related with the main contribution of one-gluon exchange in the vertex functions and is determined by the rules of \( 1/N_c \)-expansion [31,32]. For diquarks containing c-quark one common constant \( g_c = 0.857 \) is obtained. The model parameters are determined by experimental values of masses \( \Lambda_c, \Sigma_c, \Xi^{(A)}_c, \Sigma^*_c, \Xi^*_c \) [33]. The quark masses \( m_u = m_d = m = 0.495 \) GeV, \( m_s \) = 0.77 GeV and \( m_c = 1.655 \) GeV determine the region, in which it is possible to calculate the mass spectrum of all S-wave charmed baryons.

The results of calculations are shown in Table 2. In the parentheses are mentioned the experimental data [33]. In the calculations we assume the absence of the mixing of \( \Xi^{(A)}_c \) and \( \Xi^{(S)}_c \) baryons, that is in agree with result of the paper [34]. The fact is interesting, that the behaviour of Chew-Mandelstam functions which determine the dynamics of quark interactions for particles \( \Xi^{(S)}_c \), is considerably different from one for other S-wave charmed baryons. This state is excited [18] and at considered values of parameters do not arise a bound state. In order to obtain the bound excited state we introduce the quark mass shifts, which take into account the confinement potential effectively [35].

The change of quark mass values \( m_u^* = 0.41 \) GeV, \( m_s^* = 0.985 \) GeV, \( m_c^* = 1.4 \) GeV lead to the change of behaviour of Chew-Mandelstam functions and allows to receive the bound state for \( \Xi^{(S)}_c \)-baryon.

In considered calculation the masses of charmed baryons, containing two s-quarks, are overestimated in comparison with the papers [14-18]. It happens according to the choice of
the large value of s-quark mass, which allows to obtain the correct splitting of states $\Sigma_c$ and $\Xi^{(A)}_c$.

4 Conclusion

In the present paper the suggested method of the approximate solution of the relativistic three-particle equations allows us to calculate the mass spectrum of S-wave charmed baryons.

The fact, that the forces of interaction for light quarks in charmed baryons $g_1$ and $g_0$ is actually such, as well as for diquarks in the bootstrap quark model for the usual baryons. Because of the rules of $1/N_c$-expansion the diquark forces are determined by one-gluon exchange.

In the framework of suggested method it is possible to calculate the electrical formfactors and charge radii of charmed baryons, because we found the three-quark amplitudes $A_0$ and $A_1$.

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Table 1.
Coefficient of Chew-Mandelstam functions.

| $J^P$ | $\gamma(J)$ | $\beta(J)$ | $\delta(J)$ |
|-------|-------------|-------------|-------------|
| $0^+$ | $\frac{1}{2}$ | $-\frac{1}{2} \left( \frac{m_i - m_k}{m_i + m_k} \right)^2$ | 0 |
| $1^+$ | $\frac{1}{3}$ | $\frac{1}{6} \left( \frac{8m_i m_k}{(m_i + m_k)^2} - 1 \right)$ | $-\frac{1}{6} (m_i - m_k)^2$ |

Table 2.
Masses of charmed baryons of the multiplets $J^P = \frac{1}{2}^+, \frac{3}{2}^+$.  

| Quark contents | $J^P = \frac{1}{2}$ | M(GeV) | $J^P = \frac{3}{2}$ | M(GeV) |
|----------------|---------------------|--------|---------------------|--------|
| udc            | $\Lambda^+_c$       | 2.284(2.285) | ———— | ———— |
| uuc,udc,ddc    | $\Sigma^{++}_c$, $\Sigma^{++}_c$, $\Sigma^+_c$ | 2.458(2.455) | 2.516(2.519) |
| usc,dsc        | $\Xi^{+,0(A)}_c$    | 2.467(2.467) | $\Xi^{++}_c$ | 2.725(2.645) |
| usc,dsc        | $\Xi^{+,0(S)}_c$    | 2.565(2.562) | ———— | ———— |
| ssc            | $\Omega^0_c$        | 2.806(2.704) | $\Omega^0_c$ | 3.108 |
| ccu,ccd        | $\Omega^{++}_c$     | 3.527 | $\Omega^{++}_c$ | 3.597 |
| ccs            | $\Omega^+_{ccs}$    | 3.598 | $\Omega^+_c$ | 3.700 |
| ccc            | ———— | ———— | ———— | 4.792 |

The parameters of model: $\lambda_q = 10.7$ and $\lambda_c = 6.5$ are the cut-off parameter for light quarks and $c$-quark respectively, $g_0 = 0.70$ and $g_1 = 0.56$ are the vertex constant for light diquarks with $J^P = 0^+, 1^+$, $g_c = 0.857$ for heavy diquarks. Quark masses $m = 0.495$ GeV, $m_s = 0.77$ GeV, $m_c = 1.655$ GeV. The experimental values of the baryon masses are given in parentheses [33].
A Three-quark integral equations for the S-wave baryon amplitudes

\[ \text{Multiplet } J^P = \frac{3}{2}^+. \]

1. \( \Sigma^{++,+}_c \)-baryons: \( |\Sigma^+_c\rangle = |\{q \uparrow q \uparrow c \uparrow\}\rangle \), where \( q = u, d \).

\[ A_1(s, s_{12}) = \hat{\alpha} b_1(s_{12}) L_1(s_{12}) + K_1(s_{12}) [A_1^{(c)}(s, s'_{13}) + A_1^{(c)}(s, s'_{23})] \]

(A.1)

\[ A_1^{(c)}(s, s_{13}) = \hat{\alpha} b_1^{(c)}(s_{13}) L_1^{(c)}(s_{13}) + K_1^{(c)}(s_{13}) [A_1^{(s)}(s, s'_{12}) + A_1^{(s)}(s, s'_{23})] \]

(A.2)

2. \( \Xi^{+}_c \)-baryons: \( |\Xi^+_c\rangle = |\{q \uparrow s \uparrow c \uparrow\}\rangle \), where \( q = u, d \).

\[ A_1^{(s)}(s, s_{12}) = \hat{\alpha} b_1^{(s)}(s_{12}) L_1^{(s)}(s_{12}) + K_1^{(s)}(s_{12}) [A_1^{(c)}(s, s'_{13}) + A_1^{(c)}(s, s'_{23})] \]

(A.2)

\[ A_1^{(c)}(s, s_{13}) = \hat{\alpha} b_1^{(c)}(s_{13}) L_1^{(c)}(s_{13}) + K_1^{(c)}(s_{13}) [A_1^{(s)}(s, s'_{12}) + A_1^{(s)}(s, s'_{23})] \]

\[ A_1^{(sc)}(s, s_{23}) = \hat{\alpha} b_1^{(sc)}(s_{23}) L_1^{(sc)}(s_{23}) + K_1^{(sc)}(s_{23}) [A_1^{(s)}(s, s'_{12}) + A_1^{(s)}(s, s'_{13})] \]

3. \( \Omega^{0}_c \)-baryon \( |\Omega^{0}_c\rangle = |\{c \uparrow s \uparrow s \uparrow\}\rangle \). The equations are analogous to (A.1) with replacement \( A \to A^{(as)}, A^{(c)} \to A^{(sc)} \).

4. \( \Omega^{++}_c \)-baryons: \( |\Omega^{++}_c\rangle = |\{q \uparrow c \uparrow c \uparrow\}\rangle \), \( q = u, d \). The equations are analogous to (A.1) with replacement \( A \to A^{(cc)} \).

5. \( \Omega^{+}_s \)-baryon: \( |\Omega^{+}_s\rangle = |\{s \uparrow c \uparrow c \uparrow\}\rangle \). The equations are analogous to (A.1) with replacement \( A \to A^{(cc)}, A^{(c)} \to A^{(sc)} \).

6. \( \Omega^{++}_c \)-baryon: \( |\Omega^{++}_c\rangle = |\{c \uparrow c \uparrow c \uparrow\}\rangle \).

\[ A_1^{(cc)}(s, s_{12}) = \hat{\alpha} b_1^{(cc)}(s_{12}) L_1^{(cc)}(s_{12}) + K_1^{(cc)}(s_{12}) [A_1^{(cc)}(s, s'_{13}) + A_1^{(cc)}(s, s'_{23})] \]

(A.3)
Multiplet $J^P = \frac{1}{2}^+$. 

1. $\Sigma_{c}^{+, +0}$-baryons: $|\Sigma_{c}^{+}\rangle = \sqrt{\frac{1}{2}}\{|u \uparrow d \uparrow c \downarrow\} - \frac{1}{\sqrt{6}}\{|u \uparrow d \downarrow c \uparrow\} - \frac{1}{\sqrt{6}}\{|u \downarrow d \uparrow c \downarrow\}$. The spin-flavour parts of the wave functions for $\Sigma_{c}^{++}$ and $\Sigma_{c}^{0}$-baryons are derived with replacement $(ud) \rightarrow (uu)$ and $(ud) \rightarrow (dd)$ respectively.

$$A_1(s, s_{12}) = \hat{\alpha}b_1(s_{12})L_1(s_{12}) + K_1(s_{12})\left[\frac{1}{4}A_1^{(c)}(s, s_{13}) + \frac{3}{4}A_0^{(c)}(s, s'_{13}) + \frac{1}{4}A_1^{(c)}(s, s'_{23}) + \frac{3}{4}A_0^{(c)}(s, s'_{23})\right]$$

$$A_1^{(c)}(s, s_{13}) = \hat{\alpha}b_1^{(c)}(s_{13})L_1^{(c)}(s_{13}) + K_1^{(c)}(s_{13})\left[\frac{1}{2}A_1(s, s_{12}') + \frac{3}{4}A_0^{(c)}(s, s'_{12}) - \frac{1}{4}A_1^{(c)}(s, s'_{12}) + \frac{1}{2}A_1(s, s'_{23}) + \frac{3}{4}A_0^{(c)}(s, s'_{23}) - \frac{1}{4}A_1^{(c)}(s, s'_{23})\right]$$

(A.4)

$$A_0^{(c)}(s, s_{23}) = \hat{\alpha}b_0^{(c)}(s_{23})L_0^{(c)}(s_{23}) + K_0^{(c)}(s_{23})\left[\frac{1}{2}A_1(s, s_{12}') + \frac{3}{4}A_0^{(c)}(s, s'_{12}) + \frac{1}{4}A_1^{(c)}(s, s'_{12}) + \frac{1}{2}A_1(s, s'_{23}) + \frac{3}{4}A_0^{(c)}(s, s'_{23}) + \frac{1}{4}A_1^{(c)}(s, s'_{23})\right]$$

2. $\Lambda_{c}^{+}$-baryon: $|\Lambda_{c}^{+}\rangle = \frac{1}{\sqrt{2}}\{|u \uparrow d \downarrow c \uparrow\} - \frac{1}{\sqrt{2}}\{|u \downarrow d \uparrow c \downarrow\}$

$$A_0(s, s_{12}) = \hat{\alpha}b_0(s_{12})L_0(s_{12}) + K_0(s_{12})\left[\frac{1}{4}A_0^{(c)}(s, s'_{13}) + \frac{3}{4}A_1^{(c)}(s, s'_{13}) + \frac{1}{4}A_0^{(c)}(s, s'_{23}) + \frac{3}{4}A_1^{(c)}(s, s'_{23})\right]$$

$$A_1^{(c)}(s, s_{13}) = \hat{\alpha}b_1^{(c)}(s_{13})L_1^{(c)}(s_{13}) + K_1^{(c)}(s_{13})\left[\frac{1}{2}A_0(s, s_{12}') + \frac{3}{4}A_1^{(c)}(s, s'_{12}) + \frac{1}{4}A_1^{(c)}(s, s'_{12}) + \frac{1}{2}A_0(s, s'_{23}) + \frac{3}{4}A_1^{(c)}(s, s'_{23}) + \frac{1}{4}A_1^{(c)}(s, s'_{23})\right]$$

(A.5)

$$A_0^{(c)}(s, s_{23}) = \hat{\alpha}b_0^{(c)}(s_{23})L_0^{(c)}(s_{23}) + K_0^{(c)}(s_{23})\left[\frac{1}{2}A_0(s, s_{12}') + \frac{3}{4}A_1^{(c)}(s, s'_{12}) - \frac{1}{4}A_0^{(c)}(s, s'_{12}) + \frac{1}{2}A_0(s, s'_{13}) + \frac{3}{4}A_1^{(c)}(s, s'_{13}) - \frac{1}{4}A_0^{(c)}(s, s'_{13})\right]$$

3. $\Xi_{c}^{+, 0(A)}$-baryons: $|\Xi_{c}^{+, 0(A)}\rangle = \frac{1}{\sqrt{2}}\{|c \uparrow s \uparrow u \downarrow\} - \frac{1}{\sqrt{2}}\{|c \uparrow s \downarrow u \uparrow\}$ The spin-flavour part of the wave function for $\Xi_{c}^{0(A)}$-baryon are derived with replacement $(su) \rightarrow (sd)$. 

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\begin{align*}
A_0^{(s)}(s, s_{12}) &= \hat{a} b_0^{(s)}(s_{12}) L_0^{(s)}(s_{12}) + K_0^{(s)}(s_{12}) \frac{3}{8} A_1^{(c)}(s, s'_{13}) + \frac{1}{8} A_0^{(c)}(s, s'_{12}) + \frac{3}{8} A_1^{(sc)}(s, s'_{13}) + \\
&\quad + \frac{1}{8} A_0^{(sc)}(s, s'_{13}) + \frac{3}{8} A_1^{(c)}(s, s'_{23}) + \frac{1}{8} A_0^{(c)}(s, s'_{23}) + \frac{3}{8} A_1^{(sc)}(s, s'_{23}) + \frac{1}{8} A_0^{(sc)}(s, s'_{23})
\end{align*}

\begin{align*}
A_1^{(c)}(s, s_{13}) &= \hat{a} b_1^{(c)}(s_{13}) L_1^{(c)}(s_{13}) + K_1^{(c)}(s_{13}) \frac{1}{2} A_0^{(c)}(s, s'_{12}) + \frac{1}{4} A_1^{(sc)}(s, s'_{12}) + \frac{1}{4} A_0^{(sc)}(s, s'_{12}) + \\
&\quad + \frac{1}{2} A_0^{(c)}(s, s'_{23}) + \frac{1}{4} A_1^{(sc)}(s, s'_{23}) + \frac{1}{4} A_0^{(sc)}(s, s'_{23})
\end{align*}

\begin{align*}
A_0^{(c)}(s, s_{13}) &= \hat{a} b_0^{(c)}(s_{13}) L_0^{(c)}(s_{13}) + K_0^{(c)}(s_{13}) \frac{1}{2} A_0^{(c)}(s, s'_{12}) + \frac{3}{4} A_1^{(sc)}(s, s'_{12}) - \frac{1}{4} A_0^{(sc)}(s, s'_{12}) + \\
&\quad + \frac{1}{2} A_0^{(c)}(s, s'_{23}) + \frac{3}{4} A_1^{(sc)}(s, s'_{23}) - \frac{1}{4} A_0^{(sc)}(s, s'_{23})
\end{align*}

\begin{align*}
A_1^{(sc)}(s, s_{23}) &= \hat{a} b_1^{(sc)}(s_{23}) L_1^{(sc)}(s_{23}) + K_1^{(sc)}(s_{23}) \frac{1}{2} A_0^{(s)}(s, s'_{12}) + \frac{1}{4} A_1^{(c)}(s, s'_{12}) + \frac{1}{4} A_0^{(c)}(s, s'_{12}) + \\
&\quad + \frac{1}{2} A_0^{(s)}(s, s'_{13}) + \frac{1}{4} A_1^{(c)}(s, s'_{13}) + \frac{1}{4} A_0^{(c)}(s, s'_{13})
\end{align*}

\begin{align*}
A_0^{(sc)}(s, s_{23}) &= \hat{a} b_0^{(sc)}(s_{23}) L_0^{(sc)}(s_{23}) + K_0^{(sc)}(s_{23}) \frac{1}{2} A_0^{(s)}(s, s'_{12}) + \frac{3}{4} A_1^{(c)}(s, s'_{12}) - \frac{1}{4} A_0^{(c)}(s, s'_{12}) + \\
&\quad + \frac{1}{2} A_0^{(s)}(s, s'_{13}) + \frac{3}{4} A_1^{(c)}(s, s'_{13}) - \frac{1}{4} A_0^{(c)}(s, s'_{13})
\end{align*}

4. $\Xi_c^{(0)}(S)$-baryons: $|\Xi_c^{(0)}(S)\rangle = \sqrt{\frac{2}{3}} \{(u \uparrow s \uparrow c \downarrow)\} - \frac{1}{\sqrt{6}} \{(u \uparrow s \downarrow c \uparrow)\} - \frac{1}{\sqrt{6}} \{(u \downarrow s \uparrow c \uparrow)\}$. The spin-flavour part of the wave function for $\Xi_c^{(0)}(S)$-baryon are derived with replacement $(us) \rightarrow (ds)$. 

\begin{align*}
A_1^{(s)}(s, s_{12}) &= \hat{a} b_1^{(s)}(s_{12}) L_1^{(s)}(s_{12}) + K_1^{(s)}(s_{12}) \frac{3}{8} A_1^{(c)}(s, s'_{13}) + \frac{1}{8} A_0^{(c)}(s, s'_{13}) + \frac{3}{8} A_1^{(sc)}(s, s'_{13}) + \\
&\quad + \frac{3}{8} A_0^{(sc)}(s, s'_{13}) + \frac{1}{8} A_1^{(c)}(s, s'_{23}) + \frac{3}{8} A_0^{(c)}(s, s'_{23}) + \frac{3}{8} A_1^{(sc)}(s, s'_{23}) + \frac{3}{8} A_0^{(sc)}(s, s'_{23})
\end{align*}

\begin{align*}
A_1^{(c)}(s, s_{13}) &= \hat{a} b_1^{(c)}(s_{13}) L_1^{(c)}(s_{13}) + K_1^{(c)}(s_{13}) \frac{1}{2} A_1^{(s)}(s, s'_{12}) - \frac{1}{4} A_1^{(sc)}(s, s'_{12}) + \frac{3}{4} A_0^{(sc)}(s, s'_{12}) +
\end{align*}
\[+ \frac{1}{2} A_1^{(s)}(s, s'_{23}) - \frac{1}{4} A_1^{(sc)}(s, s'_{23}) + \frac{3}{4} A_0^{(sc)}(s, s'_{23})] \]

\[A_0^{(c)}(s, s_{13}) = \hat{\alpha} b_0^{(c)}(s_{13}) L_0^{(c)}(s_{13}) + K_0^{(c)}(s_{13}) \left[ \frac{1}{2} A_1^{(s)}(s, s'_{12}) + \frac{1}{4} A_1^{(sc)}(s, s'_{12}) + \frac{1}{4} A_0^{(sc)}(s, s'_{12}) + \right. \]
\[\left. + \frac{1}{2} A_1^{(s)}(s, s'_{23}) + \frac{1}{4} A_1^{(sc)}(s, s'_{23}) + \frac{1}{4} A_0^{(sc)}(s, s'_{23}) \right] \quad (A.7) \]

\[A_1^{(sc)}(s, s_{23}) = \hat{\alpha} b_1^{(sc)}(s_{23}) L_1^{(sc)}(s_{23}) + K_1^{(sc)}(s_{23}) \frac{1}{2} A_1^{(s)}(s, s'_{12}) - \frac{1}{4} A_1^{(c)}(s, s'_{12}) + \frac{3}{4} A_0^{(c)}(s, s'_{12}) + \]
\[+ \frac{1}{2} A_1^{(s)}(s, s'_{13}) - \frac{1}{4} A_1^{(c)}(s, s'_{13}) + \frac{3}{4} A_0^{(c)}(s, s'_{13})] \]

\[A_0^{(sc)}(s, s_{23}) = \hat{\alpha} b_0^{(sc)}(s_{23}) L_0^{(sc)}(s_{23}) + K_0^{(sc)}(s_{23}) \frac{1}{2} A_1^{(s)}(s, s'_{12}) + \frac{1}{4} A_1^{(c)}(s, s'_{12}) + \frac{1}{4} A_0^{(c)}(s, s'_{12}) + \]
\[+ \frac{1}{2} A_1^{(s)}(s, s'_{13}) + \frac{1}{4} A_1^{(c)}(s, s'_{13}) + \frac{1}{4} A_0^{(c)}(s, s'_{13})] \]

5. \( \Omega_c^{(0)} \)-baryon: \( |\Omega_c^{(0)}\rangle = \sqrt{\frac{2}{3}} |s \uparrow s \uparrow c \downarrow \rangle - \frac{1}{\sqrt{3}} |s \uparrow s \downarrow c \uparrow \rangle \). The equations are analogous to (A.4) with replacement \( A \to A^{(ss)}, A^{(c)} \to A^{(sc)} \).

6. \( \Omega_{ccq}^{++} \)-baryons: \( |\Omega_{ccq}^{++}\rangle = \sqrt{\frac{2}{3}} |c \uparrow c \uparrow u \downarrow \rangle - \frac{1}{\sqrt{3}} |c \uparrow c \downarrow u \uparrow \rangle \). The spin-flavour part of the wave function for \( \Omega_{ccq}^{++} \)-particle are derived with replacement \( u \to d \). The system of integral equations is analogous to (A.4) with replacement \( A \to A^{(cc)} \).

7. \( \Omega_{ccs}^{++} \)-baryon: \( |\Omega_{ccs}^{++}\rangle = \sqrt{\frac{2}{3}} |c \uparrow c \uparrow s \downarrow \rangle - \frac{1}{\sqrt{3}} |c \uparrow c \downarrow s \uparrow \rangle \). The system of integral equations is analogous to (A.4) with replacement \( A \to A^{(cc)}, A^{(c)} \to A^{(sc)} \).

**B** The systems of approximate algebraic equations for multiplets of charmed baryons \( J^P = \frac{1}{2}^+ \), \( J^P = \frac{3}{2}^+ \)

\[\text{Multiplet } J^P = \frac{3}{2}^+ \]

1. \( \Sigma_{c}^{+++},^{++,0} \)-baryons:
\[ \alpha_1(s, s_0) = \hat{\alpha} + 2\alpha_1^{(c)}(s, s_0) I_{1,1}(s, s_0) \frac{b_1^{(c)}(s_0)}{b_1(s_0)} \]

(B.1)

\[ \alpha_1^{(c)}(s, s_0) = \hat{\alpha} + \alpha_1^{(c)}(s, s_0) I_{1,1}^{(c)}(s, s_0) \frac{b_1^{(c)}(s_0)}{b_1^{(c)}(s_0)} + \alpha_1^{(c)}(s, s_0) I_{1,1}^{(c)}(s, s_0) \]

2. \(\Xi^{s+0}\)-baryons:

\[ \alpha_1^{(s)}(s, s_0) = \hat{\alpha} + \alpha_1^{(s)}(s, s_0) I_{1,1}^{(s)}(s, s_0) \frac{b_1^{(s)}(s_0)}{b_1^{(s)}(s_0)} + \alpha_1^{(s)}(s, s_0) I_{1,1}^{(s)}(s, s_0) \]

(B.2)

3. \(\Omega_{c}^{s0}\)-baryon: The equations are analogous to (A.8) with replacement \(\alpha \rightarrow \alpha^{(ss)}, \alpha^{(c)} \rightarrow \alpha^{(sc)}, b^{(c)} \rightarrow b^{(sc)}, b \rightarrow b^{(ss)}\).

4. \(\Omega_{c}^{s++}++\)-baryons: The equations are analogous to (A.8) with replacement \(\alpha \rightarrow \alpha^{(cc)}, b \rightarrow b^{(cc)}\).

5. \(\Omega_{s}^{s+}\)-baryon: The equations are analogous to (A.8) with replacement \(\alpha \rightarrow \alpha^{(cc)}, \alpha^{(c)} \rightarrow \alpha^{(sc)}, b \rightarrow b^{(cc)}, b^{(c)} \rightarrow b^{(sc)}\).

6. \(\Omega_{c}^{s+}\)-baryon:

\[ \alpha_1^{(cc)}(s, s_0) = \hat{\alpha} + \alpha_1^{(cc)}(s, s_0) 2I_{1,1}^{(cc)}(s, s_0) \]

(B.3)
Multiplet \( J^P = \frac{1}{2}^+ \).

1. \( \Sigma^+ c^{++}, c^{+0} \)-baryons:

\[
\alpha_1(s, s_0) = \hat{\alpha} + \frac{1}{2} \alpha_1^{(c)}(s, s_0) I_{1,1}(s, s_0) \frac{b_1^{(c)}(s_0)}{b_1^{(s_0)}} + \frac{3}{2} \alpha_0^{(c)}(s, s_0) I_{1,0}(s, s_0) \frac{b_0^{(c)}(s_0)}{b_1^{(s_0)}}
\]

\[
\alpha_1^{(c)}(s, s_0) = \hat{\alpha} + \alpha_1^{(c)}(s, s_0) I_{1,1}^{(c)}(s, s_0) \frac{b_{1_0}^{(c)}}{b_1^{(s_0)}} + \frac{3}{2} \alpha_0^{(c)}(s, s_0) I_{1,0}^{(c)}(s, s_0) \frac{b_0^{(c)}(s_0)}{b_1^{(s_0)}} - \frac{1}{2} \alpha_1^{(c)}(s, s_0) I_{1,1}^{(c)}(s, s_0)
\]

\[\text{(B.4)}\]

\[
\alpha_0^{(c)}(s, s_0) = \hat{\alpha} + \alpha_1^{(c)}(s, s_0) I_{0,0}^{(c)}(s, s_0) \frac{b_1^{(c)}}{b_0^{(c)}} + \frac{1}{2} \alpha_0^{(c)}(s, s_0) I_{0,0}^{(c)}(s, s_0)
\]

\[\text{(B.5)}\]

2. \( \Lambda^+ c^+ \)-baryon:

\[
\alpha_0(s, s_0) = \hat{\alpha} + \frac{1}{2} \alpha_0^{(c)}(s, s_0) I_{0,0}(s, s_0) \frac{b_0^{(c)}(s_0)}{b_0^{(s_0)}} + \frac{3}{2} \alpha_1^{(c)}(s, s_0) I_{0,1}(s, s_0) \frac{b_1^{(c)}(s_0)}{b_0^{(s_0)}}
\]

\[
\alpha_1^{(c)}(s, s_0) = \hat{\alpha} + \alpha_0^{(c)}(s, s_0) I_{1,0}^{(c)}(s, s_0) \frac{b_{1_0}^{(c)}}{b_1^{(s_0)}} + \frac{1}{2} \alpha_0^{(c)}(s, s_0) I_{1,1}^{(c)}(s, s_0) \frac{b_0^{(c)}(s_0)}{b_1^{(s_0)}}
\]

\[\text{(B.5)}\]

\[
\alpha_0^{(c)}(s, s_0) = \hat{\alpha} + \alpha_0^{(c)}(s, s_0) I_{0,0}^{(c)}(s, s_0) \frac{b_0^{(c)}}{b_0^{(s_0)}} + \frac{3}{2} \alpha_1^{(c)}(s, s_0) I_{0,1}^{(c)}(s, s_0) \frac{b_1^{(c)}}{b_0^{(s_0)}} - \frac{1}{2} \alpha_0^{(c)}(s, s_0) I_{0,0}^{(c)}(s, s_0)
\]

3. \( \Xi^+_c \)-baryons:
\[
\alpha_0^{(s)}(s, s_0) = \hat{\alpha} + \frac{3}{4} \alpha_1^{(c)}(s, s_0) I_{0,1}^{(s)}(s, s_0) \frac{b_1^{(s)}(s_0)}{b_0^{(s)}(s_0)} + \frac{1}{4} \alpha_0^{(c)}(s, s_0) I_{0,0}^{(s)}(s, s_0) \frac{b_0^{(c)}(s_0)}{b_0^{(s)}(s_0)} + \\
+ \frac{3}{4} \alpha_1^{(sc)}(s, s_0) I_{0,1}^{(s)}(s, s_0) \frac{b_1^{(sc)}(s_0)}{b_0^{(sc)}(s_0)} + \frac{1}{4} \alpha_0^{(sc)}(s, s_0) I_{0,0}^{(s)}(s, s_0) \frac{b_0^{(sc)}(s_0)}{b_0^{(sc)}(s_0)} + \\
\alpha_1^{(c)}(s, s_0) = \hat{\alpha} + \alpha_0^{(s)}(s, s_0) I_{1,0}^{(c)}(s, s_0) \frac{b_0^{(s)}(s_0)}{b_0^{(c)}(s_0)} + \\
+ \frac{1}{2} \alpha_1^{(sc)}(s, s_0) I_{1,1}^{(c)}(s, s_0) \frac{b_1^{(sc)}(s_0)}{b_1^{(sc)}(s_0)} + \frac{1}{2} \alpha_0^{(sc)}(s, s_0) I_{0,0}^{(c)}(s, s_0) \frac{b_0^{(sc)}(s_0)}{b_0^{(sc)}(s_0)} + \\
\alpha_0^{(c)}(s, s_0) = \hat{\alpha} + \alpha_0^{(s)}(s, s_0) I_{0,0}^{(c)}(s, s_0) \frac{b_0^{(s)}(s_0)}{b_0^{(c)}(s_0)} + \\
+ \frac{3}{2} \alpha_1^{(sc)}(s, s_0) I_{0,1}^{(c)}(s, s_0) \frac{b_1^{(sc)}(s_0)}{b_0^{(sc)}(s_0)} - \frac{1}{2} \alpha_0^{(c)}(s, s_0) I_{0,0}^{(c)}(s, s_0) \frac{b_0^{(c)}(s_0)}{b_0^{(c)}(s_0)}
\]

\[\text{(B.6)}\]

4. \(\Xi_{c}^{+,0(s)}\)-baryons:

\[
\alpha_1^{(s)}(s, s_0) = \hat{\alpha} + \frac{1}{4} \alpha_1^{(c)}(s, s_0) I_{1,1}^{(s)}(s, s_0) \frac{b_1^{(c)}(s_0)}{b_0^{(s)}(s_0)} + \frac{3}{4} \alpha_0^{(c)}(s, s_0) I_{1,0}^{(s)}(s, s_0) \frac{b_0^{(c)}(s_0)}{b_1^{(s)}(s_0)} + \\
+ \frac{1}{4} \alpha_1^{(sc)}(s, s_0) I_{1,1}^{(s)}(s, s_0) \frac{b_1^{(sc)}(s_0)}{b_1^{(sc)}(s_0)} + \frac{3}{4} \alpha_0^{(sc)}(s, s_0) I_{1,0}^{(s)}(s, s_0) \frac{b_0^{(sc)}(s_0)}{b_1^{(sc)}(s_0)}
\]

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\[ \alpha_1^{(c)}(s, s_0) = \hat{\alpha} + \alpha_1^{(s)}(s, s_0)I_{11}(c, s, s_0) \frac{b_1^{(s)}(s_0)}{b_1^{(c)}(s_0)} - \frac{1}{2} \alpha_1^{(sc)}(s, s_0)I_{11}(c, s, s_0) \frac{b_1^{(sc)}(s_0)}{b_1^{(c)}(s_0)} + \frac{3}{2} \alpha_0^{(sc)}(s, s_0)I_{10}(c, s, s_0) \frac{b_0^{(sc)}(s_0)}{b_1^{(c)}(s_0)} \]

\[ \alpha_0^{(c)}(s, s_0) = \hat{\alpha} + \alpha_1^{(s)}(s, s_0)I_{01}(c, s, s_0) \frac{b_1^{(s)}(s_0)}{b_0^{(c)}(s_0)} + \frac{1}{2} \alpha_1^{(sc)}(s, s_0)I_{01}(c, s, s_0) \frac{b_1^{(sc)}(s_0)}{b_0^{(c)}(s_0)} + \frac{1}{2} \alpha_0^{(sc)}(s, s_0)I_{00}(c, s, s_0) \frac{b_0^{(sc)}(s_0)}{b_0^{(c)}(s_0)} \]

(B.7)

5. \( \Omega^{0}_{cc}\)-baryon: The equations are analogous to (B.4) with replacement \( \alpha \rightarrow \alpha^{(ss)} \), \( \alpha^{(c)} \rightarrow \alpha^{(sc)} \), \( b^{(c)} \rightarrow b^{(sc)} \), \( b \rightarrow b^{(ss)} \).

6. \( \Omega^{++}_{ccq}\)-baryons: The equations are analogous to (B.4) with replacement \( \alpha \rightarrow \alpha^{(cc)} \), \( b \rightarrow b^{(cc)} \).

7. \( \Omega^{+}_{ccs}\)-baryon: The equations are analogous to (B.4) with replacement \( \alpha \rightarrow \alpha^{(cc)} \), \( \alpha^{(c)} \rightarrow \alpha^{(sc)} \), \( b \rightarrow b^{(cc)} \), \( b^{(c)} \rightarrow b^{(sc)} \).
Fig. 1. Diagrams which correspond to: a – production of three quarks, b–e – subsequent pair interaction.

Fig. 2. Combination of diagrams in accordance with the last interaction of the particles.
Fig. 3. Graphic representation of the equation for the amplitude $A_1(s, s_{12})$ (formulae (4)-(6)).

Fig. 4. The contours of integration 1, 2 in the complex plane $s_{13}$ for the function $I_{f_1, f_2}(s, s_0)$. 
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