School of Particles and Accelerators
Institute for Research in Fundamental Sciences (IPM)

First IPM meeting on

LHC Physics

April 20-24, 2009

Edited by

A. Moshaii
Tarbiat Modares University and IPM
S. Paktinat
IPM
A. Khorramian
Semnan University and IPM

IPM
December 2009

First IPM meeting on *LHC Physics*, April 20-24, 2009

was organized by

School of Particles and Accelerators, Institute for Research in Fundamental Sciences (IPM)

and sponsored by

Main Sponser: IPM

Other Sponsers:
CERN(CMS)
Center in Excellence in Physics (CEP),
Physics Department, Sharif University of Technology

International Advisory Board

F. Ardalan (Sharif University of Technology and IPM)
M. Baarmand (Florida Institute of Technology)
D. Denegri (Saclay, CERN)
J. Ellis (CERN)
L. Pape (ETH Zurich, CERN)
M. Spiropulu (CERN)
T. Virdee (Imperial College, CERN)

Organizing Committee

H. Arfaei (Sharif University of Technology and IPM)
A. Khorramian (Semnan University and IPM)
M. Mohammadi Najafabadi (IPM)
A. Moshaii (Tarbiat Modares University and IPM)
S. Paktinat (IPM)
S. Rouhani (Sharif University of Technology and IPM)

Contents

FIRST IPM MEETING
ON LHC PHYSICS
VOL. 1, NO. 1

School of Particles and Accelerators
Institute for Research in
Fundamental Sciences (IPM)
(p. 57)
Isfahan, April 20-24, 2009

6 Investigation of the D_{s1} structure via B_c to $D_{s1}l^+l^-/\nu\bar{\nu}$ transitions in QCD

M. Ghanaatian[a], R. Khosravi[b]

[a] Physics Department , Payame Noor University, Iran
[b] Physics Department , Shiraz University, Shiraz 71454, Iran

Abstract. We investigate the structure of the $D_{s1}(2460, 2536)(J^P = 1^+)$ mesons via analyzing the semileptonic $B_c \to D_{s1}l^+l^-$, $l = \tau, \mu, e$ and $B_c \to D_{s1}\nu\bar{\nu}$ transitions in the frame work of the three–point QCD sum rules. We consider the D_{s1} meson as a conventional $c\bar{s}$ meson in two ways, the pure $|c\bar{s}\rangle$ state. The obtained results for the form factors are used to evaluate the decay rates and branching ratios. Any future experimental measurement on these form factors as well as decay rates and branching fractions and their comparison with the obtained results in the present work can give considerable information about the structure of this meson.

6.1 Introduction

In this work, taking into account the gluon condensate corrections, we analyze the rare semileptonic $B_c \to D_{s1}\ l^+l^-$, $l = \tau, \mu, e$ and $B_c \to D_{s1}\nu\bar{\nu}$ transitions in three–point QCD sum rules (3PSR) approach. Note that, the $B_c \to (D^*, D_s^*, D_{s1}(2460))\nu\bar{\nu}$ transitions have been studied in Ref. [1], but assuming the D_{s1} only as $c\bar{s}$. The $B_c \to D_q\ l^+l^-/\nu\bar{\nu}$ [2], $B_c \to D_q^*\ l^+l^-$, $(q = d, s)$ [3] transitions have also been analyzed in the same framework.

The heavy B_c meson contains two heavy quarks b and c with different charges. This meson is similar to the charmonium and bottomonium in the spectroscopy, but in contrast to the charmonium and bottomonium, the B_c decays only via weak interaction and has a long lifetime. The study of the B_c transitions are useful for more precise determination of the Cabibbo, Kabayashi, Maskawa (CKM) matrix elements in the weak decays.

The rare semileptonic $B_c \to D_{s1}l^+l^-/\nu\bar{\nu}$ decays occur at loop level by electroweak penguin and weak box diagrams in the standard model (SM) via the flavor changing neutral current (FCNC) transition of $b \to sl^+l^-$. The FCNC decays of B_c meson are sensitive to new physics (NP) contributions to penguin operators. Therefore, the study of such FCNC transitions can improve the information about:

- The CP violation, T violation and polarization asymmetries in $b \to s$ penguin channels, that occur in weak interactions .
- New operators or operators that are subdominant in the SM ,
- Establishing NP and flavor physics beyond the SM.

To obtain the form factors of the semileptonic $B_c \to D_{s1}(2460[2536])$ transitions, first, we will suppose the $D_{s1}(2460)$ and $D_{s1}(2536)$ axial vector mesons as the pure $|c\bar{s}\rangle$ state and calculate the related form factors. Second, we will consider the D_{s1} meson as a mixture of two components $|D_{s1}1\rangle$ and $|D_{s1}2\rangle$ states and calculate the form factors of the $B_c \to D_{s1}1$ and $B_c \to D_{s1}2$ transitions. With the help of Eq. (6.1) and the definition of the form factors which will be presented in the next section, we will derive the transition form factors of $B_c \to D_{s1}(2460[2536])$ decays as a function of the mixing angle θ_s. The future experimental study of such rare decays and comparison of the results with the predictions of theoretical calculations can improve the information about the structure of D_{s1} meson and the mixing angle θ_s.

6.2 The form factors of $B_c \to D_{s1}$ transition in 3PSR

To calculate the form factors within three-point QCD sum rules method, the following three-point correlation functions are used [1,2,3,4]:

$$\begin{aligned}
\Pi^{V-A}_{\mu\nu}(p^2,p'^2,q^2) &= i^2 \int d^4x d^4y e^{-ipx} e^{ip'y} \langle 0 \mid T[J^{D_{s1}}_{\nu}(y) J^{V-A}_{\mu}(0) J^{B_c\dagger}(x)] \mid 0\rangle, \\
\Pi^{T-PT}_{\mu\nu}(p^2,p'^2,q^2) &= i^2 \int d^4x d^4y e^{-ipx} e^{ip'y} \langle 0 \mid T[J^{D_{s1}}_{\nu}(y) J^{T-PT}_{\mu}(0) J^{B_c\dagger}(x)] \mid 0\rangle,
\end{aligned} \tag{6.1}$$

where $J^{D_{s1}}_{\nu}(y) = \bar{c}\gamma_\nu\gamma_5 s$ and $J^{B_c}(x) = \bar{c}\gamma_5 b$ are the interpolating currents of the initial and final meson states, respectively. $J^{V-A}_{\mu} = \bar{s}\gamma_\mu(1-\gamma_5)b$ and $J^{T-PT}_{\mu} = \bar{s}\,\sigma_{\mu\nu}q^{\nu}(1+\gamma_5)b$ are the vector-axial vector and tensor-pseudo tensor parts of the transition currents. In QCD sum rules approach, we can obtain the correlation function of Eq. (6.1) in two sides. The phenomenological or physical part is calculated saturating the correlator by a tower of hadrons with the same quantum numbers as interpolating currents. The QCD or theoretical part, on the other side is obtained in terms of the quarks and gluons interacting in the QCD vacuum. To drive the phenomenological part of the correlators given in Eq. (6.1), two complete sets of intermediate states with the same quantum numbers as the currents $J_{D_{s1}}$ and J_{B_c} are inserted. This procedure leads to the following representations of the above-mentioned correlators:

$$\begin{aligned}
\Pi^{V-A}_{\mu\nu}(p^2,p'^2,q^2) &= \frac{\langle 0 \mid J^{D_{s1}}_{\nu} \mid D_{s1}(p',\varepsilon)\rangle \langle D_{s1}(p',\varepsilon) \mid J^{V-A}_{\mu} \mid B_c(p)\rangle \langle B_c(p) \mid J^{B_c\dagger} \mid 0\rangle}{(p'^2 - m^2_{D_{s1}})(p^2 - m^2_{B_c})} \\
&\quad + \text{higher resonances and continuum states}\,, \\
\Pi^{T-PT}_{\mu\nu}(p^2,p'^2,q^2) &= \frac{\langle 0 \mid J^{D_{s1}}_{\nu} \mid D_{s1}(p',\varepsilon)\rangle \langle D_{s1}(p',\varepsilon) \mid J^{T-PT}_{\mu} \mid B_c(p)\rangle \langle B_c(p) \mid J^{B_c\dagger} \mid 0\rangle}{(p'^2 - m^2_{D_{s1}})(p^2 - m^2_{B_c})} \\
&\quad + \text{higher resonances and continuum states}\,.
\end{aligned} \tag{6.2}$$

The following matrix elements are defined in the standard way in terms of the leptonic decay constants of the D_{s1} and B_c mesons as:

$$\langle 0 \mid J^{\nu}_{D_{s1}} \mid D_{s1}(p',\varepsilon)\rangle = f_{D_{s1}} m_{D_{s1}} \varepsilon^{\nu}\ ,\ \ \langle 0 \mid J_{B_c} \mid B_c(p)\rangle = i\frac{f_{B_c} m^2_{B_c}}{m_b + m_c}. \tag{6.3}$$

To parameterized the matrix elements in terms of the transition form factors considering the Lorentz invariance and parity considerations.

$$\begin{aligned}
\langle D_{s1}(p',\varepsilon) \mid \bar{s}\gamma_{\mu}\gamma_5 b \mid B_c(p)\rangle &= \frac{2A_V^{B_c\to D_{s1}}(q^2)}{m_{B_c}+m_{D_{s1}}}\ \varepsilon_{\mu\nu\alpha\beta}\varepsilon^{*\nu}p^{\alpha}p'^{\beta}, \\
\langle D_{s1}(p',\varepsilon) \mid \bar{s}\gamma_{\mu} b \mid B_c(p)\rangle &= -\,iA_0^{B_c\to D_{s1}}(q^2)(m_{B_c}+m_{D_{s1}})\varepsilon^*_{\mu} \\
&\quad + i\frac{A_1^{B_c\to D_{s1}}(q^2)}{m_{B_c}+m_{D_{s1}}}(\varepsilon^* p)P_{\mu} \\
&\quad + i\frac{A_2^{B_c\to D_{s1}}(q^2)}{m_{B_c}+m_{D_{s1}}}(\varepsilon^* p)q_{\mu}\ , \\
\langle D_{s1}(p',\varepsilon) \mid \bar{s}\sigma_{\mu\nu}q^{\nu}\gamma_5 b \mid B_c(p)\rangle &= 2\ T_V^{B_c\to D_{s1}}(q^2)\ i\varepsilon_{\mu\nu\alpha\beta}\varepsilon^{*\nu}p^{\alpha}p'^{\beta}, \\
\langle D_{s1}(p',\varepsilon) \mid \bar{s}\sigma_{\mu\nu}q^{\nu} b \mid B_c(p)\rangle &= T_0^{B_c\to D_{s1}}(q^2)[\varepsilon^*_{\mu}(m^2_{B_c}-m^2_{D_{s1}})-(\varepsilon^* p)P_{\mu}] \\
&\quad + T_1^{B_c\to D_{s1}}(q^2)\ (\varepsilon^* p)[q_{\mu}-\frac{q^2}{m^2_{B_c}-m^2_{D_{s1}}}P_{\mu}],
\end{aligned} \tag{6.4}$$

where $A_i^{B_c\to D_{s1}}(q^2)$, $i = V, 0, 1, 2$ and $T_j^{B_c\to D_{s1}}(q^2)$, $j = V, 0, 1$ are the transition form factors, $P_{\mu} = (p+p')_{\mu}$ and $q_{\mu} = (p-p')_{\mu}$. Here, q^2 is the momentum transfer squared of the Z boson (photon). In order to our calculations be simple, the following redefinitions of the transition form factors are considered :

$$\begin{aligned}
A_V'^{B_c\to D_{s1}}(q^2) &= \frac{2A_V^{B_c\to D_{s1}}(q^2)}{m_{B_c}+m_{D_{s1}}}\ , \\
A_0'^{B_c\to D_{s1}}(q^2) &= A_0^{B_c\to D_{s1}}(q^2)(m_{B_c}+m_{D_{s1}})\ , \\
A_1'^{B_c\to D_{s1}}(q^2) &= -\frac{A_1^{B_c\to D_{s1}}(q^2)}{m_{B_c}+m_{D_{s1}}}\ , \\
A_2'^{B_c\to D_{s1}}(q^2) &= -\frac{A_2^{B_c\to D_{s1}}(q^2)}{m_{B_c}+m_{D_{s1}}}\ , \\
T_V'^{B_c\to D_{s1}}(q^2) &= -2T_V^{B_c\to D_{s1}}(q^2)\ , \\
T_0'^{B_c\to D_{s1}}(q^2) &= -\ T_0^{B_c\to D_{s1}}(q^2)(m^2_{B_c}-m^2_{D_{s1}})\ , \\
T_1'^{B_c\to D_{s1}}(q^2) &= -\ T_1^{B_c\to D_{s1}}(q^2)\ .
\end{aligned} \tag{6.5}$$

Using Eq. (6.4), Eq. (6.5) and Eq. (6.3) in Eq. (6.2) and performing summation over the polarization of D_{s1} meson we obtain:

$$\begin{aligned}
\Pi_{\mu\nu}^{V-A}(p^2,p'^2,q^2) &= -\frac{f_{B_c}m_{B_c}^2}{(m_b+m_c)}\frac{f_{D_{s1}}m_{D_{s1}}}{(p'^2-m_{D_{s1}}^2)(p^2-m_{B_c}^2)} \times \Big[iA_V^{'B_c\to D_{s1}}(q^2)\varepsilon_{\mu\nu\alpha\beta}p^\alpha p'^\beta \\
&+ A_0^{'B_c\to D_{s1}}(q^2)g_{\mu\nu} + A_1^{'B_c\to D_{s1}}(q^2)P_\mu p_\nu + A_2^{'B_c\to D_{s1}}(q^2)q_\mu p_\nu\Big] \\
&+ \text{excited states}, \\
\Pi_{\mu\nu}^{T-PT}(p^2,p'^2,q^2) &= -\frac{f_{B_c}m_{B_c}^2}{(m_b+m_c)}\frac{f_{D_{s1}}m_{D_{s1}}}{(p'^2-m_{D_{s1}}^2)(p^2-m_{B_c}^2)} \times \Big[T_V^{'B_c\to D_{s1}}(q^2)\varepsilon_{\mu\nu\alpha\beta}p^\alpha p'^\beta \\
&- i\,T_0^{'B_c\to D_{s1}}(q^2)g_{\mu\nu} - i\,T_1^{'B_c\to D_{s1}}(q^2)q_\mu p_\nu\Big] + \text{excited states.}
\end{aligned} \tag{6.6}$$

To calculate the form factors, A_V', A_0', A_1', A_2', T_V', T_0' and T_1', we will choose the structures, $i\varepsilon_{\mu\nu\alpha\beta}p^\alpha p'^\beta$, $g_{\mu\nu}$, $P_\mu p_\nu$, $q_\mu p_\nu$, from $\Pi_{\mu\nu}^{V-A}$ and $\varepsilon_{\mu\nu\alpha\beta}p^\alpha p'^\beta$, $ig_{\mu\nu}$ and $iq_\mu p_\nu$ from $\Pi_{\mu\nu}^{T-PT}$, respectively.

On the QCD side, using the operator product expansion (OPE), we can obtain the correlation function in quark-gluon language in the deep Euclidean region where $p^2 \ll (m_b+m_c)^2$, $p'^2 \ll (m_c^2+m_s^2)$. For this aim, the correlators are written as:

$$\begin{aligned}
\Pi_{\mu\nu}^{V-A}(p^2,p'^2,q^2) &= i\,\Pi_V^{V-A}\varepsilon_{\mu\nu\alpha\beta}p^\alpha p'^\beta + \Pi_0^{V-A}g_{\mu\nu} + \Pi_1^{V-A}P_\mu p_\nu + \Pi_2^{V-A}q_\mu p_\nu, \\
\Pi_{\mu\nu}^{T-PT}(p^2,p'^2,q^2) &= \Pi_V^{T-PT}\varepsilon_{\mu\nu\alpha\beta}p^\alpha p'^\beta - i\,\Pi_0^{T-PT}g_{\mu\nu} - i\,\Pi_1^{T-PT}q_\mu p_\nu,
\end{aligned} \tag{6.7}$$

where, each Π_i function is defined in terms of the perturbative and non-perturbative parts as:

$$\Pi_i(p^2,p'^2,q^2) = \Pi_i^{per}(p^2,p'^2,q^2) + \Pi_i^{nonper}(p^2,p'^2,q^2)\,. \tag{6.8}$$

Performing the double Borel transformations over the variables p^2 and p'^2 on the physical as well as perturbative parts of the correlation functions and equating the coefficients of the selected structures from both sides, the sum rules for the form factors $A_i^{'B_c\to D_{s1}}$ are obtained:

$$\begin{aligned}
A_i^{'B_c\to D_{s1}} &= -\frac{(m_b+m_c)}{f_{B_c}m_{B_c}^2 f_{D_{s1}}m_{D_{s1}}} e^{\frac{m_{B_c}^2}{M_1^2}} e^{\frac{m_{D_{s1}}^2}{M_2^2}} \Bigg\{ -\frac{1}{4\pi^2}\int_{m_c^2}^{s_0'} ds' \int_{s_L}^{s_0} \rho_i^{V-A}(s,s',q^2) \\
&\quad e^{\frac{-s}{M_1^2}} e^{\frac{-s'}{M_2^2}} - iM_1^2M_2^2\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle\frac{C_i^{V-A}}{6}\Bigg\},
\end{aligned} \tag{6.9}$$

where $i = V, 0, 1, 2$ and for form factors $T_j^{'B_c\to D_{s1}}$, we get

$$\begin{aligned}
T_j^{'B_c\to D_{s1}} &= -\frac{(m_b+m_c)}{f_{B_c}m_{B_c}^2 f_{D_{s1}}m_{D_{s1}}} e^{\frac{m_{B_c}^2}{M_1^2}} e^{\frac{m_{D_{s1}}^2}{M_2^2}} \Bigg\{ -\frac{1}{4\pi^2}\int_{m_c^2}^{s_0'} ds' \int_{s_L}^{s_0} \rho_j^{T-PT}(s,s',q^2) \\
&\quad e^{\frac{-s}{M_1^2}} e^{\frac{-s'}{M_2^2}} - iM_1^2M_2^2\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle\frac{C_j^{T-PT}}{6}\Bigg\}.
\end{aligned} \tag{6.10}$$

where $j = V, 0, 1$. The s_0 and $s_0^{'}$ are the continuum thresholds in B_c and D_{s1} channels, respectively and lower bound s_L in the integrals. We calculated the explicit expressions of the coefficients $C_{i(j)}^{V-A(T-PT)}$ correspond to gluon condensates.

6.3 Numerical analysis

In this section, we present our numerical analysis of the form factors A_i ,$(i = V, 0, 1, 2)$ and T_j ,$(j = V, 0, 1)$. From the sum rules expressions of the form factors, it is clear that the main input parameters entering the expressions are gluon condensates, elements of the CKM matrix V_{tb} and V_{ts}, leptonic decay constants f_{B_c}, $f_{D_{s1}}$, $f_{D_{s1}1}$ and $f_{D_{s1}2}$, Borel parameters M_1^2 and M_2^2 as well as the continuum thresholds s_0 and s_0'. We choose the values of the condensates (at a fixed renormalization scale of about 1 GeV), leptonic decay constants , CKM matrix elements, quark and meson masses [5,6,7,8,9,10,11,12,13].

First, we would like to consider the D_{s1} meson as the pure $|c\bar{s}\rangle$ state. To calculate the branching ratios of the $B_c \to D_{s1}(2460[2536])l^+l^-/\nu\bar{\nu}$ decays, we use the total mean life time $\tau_{B_c} = (0.46 \pm 0.07)$ ps [13]. Our numerical analysis shows that the contribution of the non-perturbative part (the gluon condensate diagrams) is about 12% of the total and the main contribution comes from the perturbative part of the form factors. The values for the branching ratio of these decays are obtained as presented in Table 6.1, when only the short distance (SD) effects are considered.

MODS	BR	MODS	BR
$B_c \to D_{s1}(2460)\nu\bar{\nu}$	$(3.26 \pm 1.10) \times 10^{-7}$	$B_c \to D_{s1}(2536)\nu\bar{\nu}$	$(2.76 \pm 0.88) \times 10^{-7}$
$B_c \to D_{s1}(2460)e^+e^-$	$(5.40 \pm 1.70) \times 10^{-6}$	$B_c \to D_{s1}(2536)e^+e^-$	$(2.91 \pm 0.93) \times 10^{-6}$
$B_c \to D_{s1}(2460)\mu^+\mu^-$	$(2.27 \pm 0.95) \times 10^{-6}$	$B_c \to D_{s1}(2536)\mu^+\mu^-$	$(1.96 \pm 0.63) \times 10^{-6}$
$B_c \to D_{s1}(2460)\tau^+\tau^-$	$(1.42 \pm 0.45) \times 10^{-8}$	$B_c \to D_{s1}(2536)\tau^+\tau^-$	$(0.68 \pm 0.21) \times 10^{-8}$

Table 6.1. The branching ratios of the semileptonic $B_c \to D_{s1}(2460)l^+l^-/\nu\bar{\nu}$ and $B_c \to D_{s1}(2536)l^+l^-/\nu\bar{\nu}$ decays with SD effects.

References

1. K. Azizi, R. Khosravi, V. Bashiry, Eur. Phys. J. C 56, 357, (2008).
2. K. Azizi, R. Khosravi, Phys. Rev. D 78, 036005 (2008).
3. K. Azizi, F. Falahati, V. Bashiry, S. M. Zebarjad, Phys. Rev. D 77, 114024 (2008).
4. N. Ghahramany, R. Khosravi, K. Azizi, Phys. Rev. D 78, 116009 (2008).
5. M. A. Shifman, A. I. Vainshtein, V. I. Zakharov, Nucl. Phys. B 147, 385 (1979).

6. A. Ceccucci, Z. Ligeti,Y. Sakai, PDG, J. Phys. G 33, 139 (2006).
7. A. J. Buras, M. Muenz, Phys. Rev. D 52, 186 (1995).
8. V. Bashiry, K. Azizi, JHEP 0707, 064 (2007).
9. S. Veseli, I. Dunietz, Phys. Rev. D 54, 6803 (1996).
10. P. Colangelo, G. Nardulli, N. Paver, Z. Phys. C 57, 43 (1993).
11. V. V. Kiselev, A. V. Tkabladze, Phys. Rev. D 48, 5208 (1993).
12. T. M. Aliev, O. Yilmaz, Nuovo Cimento. A 105, 827 (1992).
13. C. Amsler et al., Particle Data Group, Phys. Lett. B 667, 1 (2008).

FIRST IPM MEETING
ON LHC PHYSICS
VOL. 1, NO. 1

School of Particles and Accelerators
Institute for Research in
Fundamental Sciences (IPM)
(p. 63)
Isfahan, April 20-24, 2009

7 Quark-Gluon Plasma Model and Origin of Magic Numbers

M. Ghanaatian[a], N.Ghahramany[b]

[a] Physics Department, Payame Noor University, Iran
[b] Physics Department, Shiraz University, Shiraz 71454, Iran

Abstract. Using Boltzman distribution in a quark-gluon plasma sample it is possible to obtain all existing magic numbers and their extensions without applying the spin and spin-orbit couplings. In this model it is assumed that in a quark-gluon thermodynamic plasma, quarks have no interactions and they are trying to form nucleons. Considering a lattice for a central quark and the surrounding quarks, using a statistical approach to find the maximum number of microstates, the origin of magic numbers is explained and a new magic number is obtained.

7.1 Introduction

There are certain elements in the universe with relative high stability and abundance whose neutron or proton numbers are called magic numbers. Historically, their stability and excited energies were first found [1] but the origin of magic numbers remained as a mystery. Of course there are some explanations about these numbers from shell model of nuclei, mainly from observation of apparent similarities between these magic numbers and nucleon numbers that fill the nuclear shells. In nuclear shell model it is assumed that the nucleon is orbiting in a nuclear spherical potential well and the energy gaps between the spectral lines obtained from such potential well correspond to the stability of nuclei. Such correspondence was not accurate enough, therefore several researchers, mainly Maria Goepert Mayer [2] included the effect of spin and spin-orbit coupling in the nuclear Hamiltonian as a perturbation from which new energy gaps were observed in more agreement with the observed magic numbers. Eventually the shell model was built explaining the nuclear structure and position of constituent particles with no satisfactory explanation about the origin of the magic numbers.

In this research it is intended to investigate the origin of these magic numbers via quark-gluon plasma media. In this model it is assumed that in a quark-gluon thermodynamic plasma in which the quark have no interaction, quark are trying to form nucleons. If we accept that the stability of a thermodynamical system is obtained when the system is in maximum disorder or maximum number of complexions, then by considering different isolated system containing one central quark and 2, 3, 4, 5, 6, 7 and at most 8 surrounding quarks embracing the central quark, these can find maximum number of microstates that correspond exactly to magic numbers. From statistical point of view it will be clear that why it is unlikely to have a central quark with 8 surrounding quarks and therefore very likely to have a central one with 2 or 3 quarks around it.

7.2 Colored quark-gluon plasma and magic numbers

The hot quark-gluon plasma (QGP) exists right after big-bang and by relativistic expansion cool down and change to proton and neutron. In the continued process of expansion, different nuclei are formed (nucleation) via Boltzman equilibrium process [3,4,5,6,7]. These formed nuclei are most stable at magic numbers. We intend to investigate how and under what conditions the quarks with color and flavor form proton and neutron and what the origin of magic numbers is in such QGP. It is not intended to describe the quark distribution after the formation of proton which is given in terms of structure functions. If the QGP is considered as a thermodynamical media then it should proceed toward maximum disorder. It should be investigated that how such system approach equilibrium. The thermodynamical state is a stable system with maximum probability state, i.e., the most probable state with maximum number of complexion.

Now consider a thermodynamical state of quarks in motion. In such QGP soup the quarks are not absolutely free. This is known from lattice QCD theory [8]. In fact in such QGP soup, the gluons connect to the nearby quarks with a force much weaker than the binding force. The QGP media is assumed as an ideal gas model. In such model consider a quark to be trying to form a nucleon capturing two quarks of different flavor. In such competing space between quarks different nucleon formation cases happen.

If there is only two d-quark with no color neighboring the central u-quark, then there is no competition and state ud_1d_2 is formed. But from standard model each quark have 3 color and d_1 and d_2 must be of different color say green and blue, therefore two competing cases exist namely, a red u with blue d_1 and green d_2 or with green d_1 and blue d_2. So we get number 2 as the first magic number.

Now consider that there are three d-quarks neighboring the central u-quark, then there are three cases namely, ud_1d_2, ud_1d_3 and ud_2d_3. If their color is also taken into account, then there are six cases in addition to the previous two cases and we get eight cases to form a proton, i.e., the second magic number. Lets consider four neighboring d-quarks. If only two of them compete, we have 2 cases and if three of them participate in this competition then we have six cases and if all four compete then we get 12 cases. The total number of cases is therefore 2+6+12=20. This is the third magic number.

It is interesting to note that in both the spin and spin-orbit coupling interpretations to explain magic numbers [9] and numerical explanation of Bagge [10], two separate series were introduced namely,(2,8,20,40,70,112) and (2,6,14,28,50,82 ,126). After the magic number 20, there was a jump from the first series to the second series to obtain number 28.

If we consider 5 quarks neighboring the central one, the competition between these 5 quarks in addition to previous ones given us 40 cases which is exactly the fourth number from the first series and for 6 and 7 neighboring quarks 70 and

112 cases are obtained and the same historical problem do exist. To resolve this problem in our model, the QGP condition is utilized and "imposed quarks" are introduced. From lattice QCD theory it is clear that quarks are not free and there exist some weak attractive forces between quarks in QGP soup and that is why it is called a soup. Now suppose there are four d quarks neighboring the central u-quark one as obtained before there are 20 cases competing to form a proton. If each d-quark is considered to be close to its own neighbors. then if for example d is absorbed by it then the closest quark to d which has the strongest attraction force to d will accompany it and participate. Lets call it d' and this is named as imposed quark.

This quark come from the second level, therefore with one imposed quark for each initial 4 d-quarks we have $ud_1d'_1$,$ud_2d'_2$, $ud_3d'_3$ and $ud_4d'_4$ and considering their colors there are 8 cases in addition to the pervious 20 cases, adding to 28 competing states to form a proton. Lets consider five d-quarks surrounding the central u-quark, then we have twenty new cases in addition to ten imposed cases adding to 50 cases 20+20+10=50, which consist of ud_1d_2, ud_1d_3, ud_1d_4, ud_1d_5, ud_2d_3, ud_2d_4, ud_2d_5, ud_3d_4, ud_3d_5, ud_4d_5 and $ud_1d'_1$, $ud_2d'_2$, $ud_3d'_3$, $ud_4d'_4$ and $ud_5d'_5$. Now lets consider six d-quarks around the central u-quark, then we have: 40+30+12=82. For seven d-quarks participating: 70+42+14=126. Eventually for eight d-quarks we get 184 cases. This is a new magic number that is obtained in this model. For more than 8 quarks one has to consider the imposed quarks from the third level and no additional magic number is obtained.

7.3 Conlusion

A quark-gluon plasma soup model is presented based upon Boltzman distribution and an alternative approach is suggested to obtain not only the existing magic numbers exactly but new magic number is introduced. A quark shell structure in the form of cubic lattice is considered to find the most probable cases and maximum number of ways that a stable element is formed corresponding to the magic numbers. While work is in progress to understand how exactly the same magic numbers appear in nuclear formation, this paper is intended to provide insight and extended the concept of magic numbers from nuclei to nucleon formation independently.

References

1. W. Elsasser, J. Phys. Radium, 4, 549 (1933).
2. M. G. Mayer, Phys. Rev., Vol. 74, 235 (1948); M. G. Mayer, Phys. Rev., Vol. 75, 1969 (1949); M. G. Mayer, and J. H. D. Jensen, Elementary Theory of Nuclear Shell Structure, John Wiley & Sons, New York, 1955.
3. A. Lefebvre et al, Nuclear Physics, A621, 199 (1997).
4. R. Kippenhahn, and A. Weigert, Stellar Structure and Evolution, Springer, Heidelberg, 1990.
5. F. Osman, N. Ghahramany, and H. Hora, Laser and Particle Beams, 23, 461-466, (2005).
6. H. Hora, Plasma Model for Surface Tension of Nuclei and the Phase Transition to the Quark Plasma, Report CERN-PS/DL-Note-91/05, 1991.

7. T. Ranscher, J. H. Applegate, J. J. Cowan, F. K. Thielmann, and M. Weiseher, Astrophys. J., 429, 499 (1994).
8. K.G. Wilson, Phys. Rev. D 10, 2445 (1974); K. G. Wilson, rev. Mod. Phys. 55, 583 (1983); K. G. Wilson et al, Phys. Rev. D49, 6720 (1994).
9. O. Haxel, J. H. D. Jensen and H. E. Suess, Zeitscher. f. Physik 128, 295 (1950).
10. E. Bagge, Naturwissenschaften 35, 376 (1948).