Heavy baryons in the large $N_c$ limit

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It is shown that in the large $N_c$ limit heavy baryon masses can be estimated quantitatively in a $1/N_c$ expansion using the Hartree approximation. The results are compared with available lattice calculations for different values of the ratio between the square root of the string tension and the heavy quark mass $\sqrt{\sigma}/m_Q$. These estimates implement important $1/N_c$ corrections and assume a string tension independent of $N_c$. Using a potential adjusted to agree with the one obtained in lattice QCD, a variational analysis of the ground state spin averaged baryon mass is performed using Gaussian Hartree wave functions. Relativistic corrections through the quark kinetic energy are included. The results provide good estimates for the first sub-leading in $1/N_c$ corrections.

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1. Introduction

QCD in the large $N_c$ limit becomes a non-trivial theory in terms of an arbitrary and fixed 't Hooft coupling $\lambda = \alpha_s N_c$ [1]. In that limit, baryons [2], unlike mesons, remain as complicated structures (for a recent review see e.g. [3,4] and references therein). This is the result of the strong coupling of mesons to baryons $O(\sqrt{N_c})$, giving baryons a light meson cloud which contributes to its mass at leading order in $N_c$. In the world of QCD with only heavy quarks, the meson cloud becomes suppressed in $\Delta_{QCD}/m_Q$, $m_Q$ being the heavy quark mass, and baryonic states become amenable to a treatment based on non-relativistic QCD. Thus, heavy baryons are a good laboratory to study the $1/N_c$ expansion. This simpler setting of QCD permits a straightforward application of the mean field approach, which will be used in the present work and which should provide a good description of baryons in the large $N_c$ and large quark mass limits.

The quantitative understanding of the $1/N_c$ expansion has become possible in the light meson sector [5], where meson masses have been determined in lattice QCD (LQCD) calculations at different values of $N_c$ and in the quenched approximation, where the leading $O(1/N_c)$ corrections are absent, and moderate $N_c$ values allow for a safe extrapolation to the large $N_c$ limit. In addition, estimates based on short distance constraints provide an analytical understanding of those results [6]. More recently, LQCD calculations of low lying baryon masses for $N_c = 3$, 5 and 7 [7,8] have opened the door for a quantitative test of the $1/N_c$ expansion in baryons as well. Those pioneering calculations, which are in the quenched approximation, have quark masses in the light to moderately heavy range. The present work is largely motivated by the possibility that such LQCD calculations could be extended to heavier quark masses, where the framework presented here would become realistically applicable.

In his seminal paper, Witten [2] discussed specifically heavy baryons in the large $N_c$ limit and invoked the mean field Hartree approximation. For heavy quarks, it is built from the simple two-body Hamiltonian, where the interaction is the OGE (one gluon exchange) (see [9] for details) for the short range part of the interaction. In addition, there are the long range confining forces, whose effects become suppressed as $m_Q$ grows, and also short distance radiative corrections must be taken into account (running of $\alpha_s$) (see [10]). Furthermore, the effects of three-body interactions are of potential interest; for a recent discussion in the quark model see Ref. [11]. They will be discussed briefly in this work.

At leading order in the $1/N_c$ expansion, the ground state of the heavy baryon will be described by a wave function which is the direct product of single-quark wave functions. Since the hyperfine interactions have spin-flavor non-singlet effects which are $O(1/N_c)$, it is clear that at leading order the ground state baryon is in the totally symmetric spin-flavor state, and the baryon has a spin-flavor contracted symmetry [12,13], which holds in the limit $N_c \rightarrow \infty$ at fixed quark mass. The effect of removing the center of mass (CM) motion is sub-leading in $1/N_c$, and can be implemented...
using standard techniques such as the Peierls–Yoccoz projection (for a review see e.g. [14,15] and references therein).

The mean field for heavy quarks at large $N_c$ has been studied in Refs. [16,17] along with possible implications for baryonic matter. This work builds on that one and compares to recent lattice calculations for $N_c = 3, 5, 7$ [7,8] after including some important $1/N_c$ effects such as the CM correction. Brief discussions of the role of hyperfine splittings as well as the expected corrections of many-body forces are also given. A previous large $N_c$ analysis has been conducted in Ref. [18].

Note that in order to have low lying baryons with different spins it is necessary to have more than one flavor of heavy quark. The mass of the baryon will then have an $O(1/N_c)$ hyperfine contribution (dependent on the spin $S$ of the baryon). The masses of ground state baryons take the form of a rotational band,

$$M_B(S) = N_c m_0 + \frac{C_{HF}}{N_c} (S(S+1) - \frac{3}{4} N_c) + O\left(\frac{1}{N_c^2}\right),$$

where $m_0$ and $C_{HF}$ are $O(1/N_c^0)$ and have an expansion in $1/N_c$, and depend on the quark mass $m_Q$. The hyperfine independent component of the baryon mass given by $m_0$ is obtained by the following combination of baryon masses:

$$m_0 = \frac{2}{N_c^2(N_c+1)(N_c+3)^2} \times \sum_{S=\frac{1}{2}}^{N_c} \left(3 + N_c(3N_c+2) - 8(N_c-3)S\right) M_B(S).$$

The baryon masses studied here will be the ones with the hyperfine effects removed, i.e., $M_B \equiv N_c m_0$. These will be later compared with the available LQCD results of Refs. [7,8,18].

Of course, for any different value of $N_c$ one has a different theory. Thus, in order to relate them one must assume that some observables are $N_c$ independent. Actually, on general grounds one has that:

$$m_0 \sqrt{\sigma} = F(N_c) \left(\frac{m_Q}{\sqrt{\sigma}}\right),$$

where $\sigma$ sets the scale of QCD and can be identified for instance with the string tension, and $m_Q$ is the heavy quark mass. $F$ is a universal function $O(N_c^0)$ which admits an expansion in $1/N_c$, and which for large $m_0$ can be more conveniently expressed as $F(N_c, m_0/\sqrt{\sigma}) = m_0/\sqrt{\sigma} f(N_c, m_0/\sqrt{\sigma})$.

The present work goes beyond Refs. [16,17] by analyzing the main $1/N_c$ contributions such as the CM effect, and relativistic corrections, and actually compares to available LQCD results. For $N_c = 3$, triply heavy baryons have been studied on the lattice as a $\Omega_{hhh}$ state [19], and also re-addressed in quark models within several schemes [10,11,20] which, however, have not addressed larger $N_c$ values.

One important goal on the lattice has been to make the quarks light as possible. Actually, quarkonium studies based on LQCD proceed always through the determination of the $Q/Q$ potential, and a subsequent solution of the non-relativistic Schrödinger equation (see e.g. [21]). The present work takes a similar point of view as a $N_c$-body problem. It should be emphasized that studying heavy baryons at varying values of $N_c$ will help with the understanding of the $1/N_c$ expansion in a setting where an analytic approach with small model dependencies can be applied.

2. Color singlet states

The starting point is the Hamiltonian for heavy quarks. Using non-relativistic heavy quark field operators $Q(x)$, the Hamiltonian is given by:

$$H = \int d^3 x \left[ -\frac{1}{2m_Q} Q^\dagger(x) \nabla^2 Q(x) + m_Q \left(\frac{1}{2} Q(x)^\dagger Q(x) \right) \right] + \frac{1}{2} \int d^3 x \, d^3 x' \, Q(x) \lambda^a(x) \frac{\lambda^a}{2} Q(x') V(x-x'),$$

(4)

where $\lambda^a$ are the $SU(N_c)$ generators in the fundamental representation, and in perturbation theory $V(r) = \alpha / r$ is the OGE interaction. Here, only two-body interactions are included. The role of many body interactions is commented below. An equivalent representation for the case of a heavy baryon is the Hamiltonian

$$H = \sum_i \left[ m_Q + \frac{p_i^2}{2m_Q} \right] + \frac{1}{4} \sum_{i<j} \lambda_a(i) \otimes \lambda^a(j) \bar{V}(x_i-x_j),$$

(5)

The $\lambda \otimes \lambda$ interaction implies exact Casimir scaling of the potential energy. Casimir scaling for the $Q\bar{Q}$ potential holds perturbatively up to two loops (there are three-loop violations) [22] and numerically on the lattice [23].

For a color singlet state the wave function is completely symmetric in the orbital and spin-flavor quantum numbers, and the baryon behaves effectively as a bosonic system. In particular, for ground state baryons the wave function is the product of a symmetric spacial wave function and a symmetric spin-flavor wave function and reads as follows:

$$\Psi(x_1, \ldots, x_N) = \psi(x_1, \ldots, x_N) \chi_{SF},$$

(6)

where $\chi_{SF}$ is the spin-flavor wave function. For excited baryon states, spin-flavor and spatial mixed symmetry states also occur. The color matrix elements for arbitrary $N_c$ in the ground state can be computed as follows. Starting with the quadratic Casimir operator for the fundamental representation given by $(F^a = \lambda^a/2)$

$$\bar{F} \cdot \bar{F} = 2 \bar{F} \cdot \bar{F} = \frac{N_c^2 - 1}{2N_c},$$

(7)

for a baryon (color singlet) state one obtains:

$$0 = \langle B | \sum_{i=1}^{N_c} |\bar{F}_i|^2 \rangle |B\rangle$$

$$= \langle B | \sum_{i=1}^{N_c} |\bar{F}_i|^2 \rangle |B\rangle + 2 \sum_{i<j} \langle B | \bar{F}_i \cdot \bar{F}_j | B \rangle$$

$$= N_c (\langle B | (\bar{F}_Q^2) | B \rangle + N_c(N_c-1) \langle B | \bar{F}_Q \cdot \bar{F}_Q | B \rangle),$$

(8)

and likewise for a meson state one obtains:

$$0 = \langle M | (\bar{F}_Q + \bar{F}_Q)^2 | M \rangle$$

$$= 2 \langle M | (\bar{F}_Q)^2 | M \rangle + 2 \langle M | \bar{F}_Q \cdot \bar{F}_Q | M \rangle$$

(9)

These equations lead to

$$\langle B | \bar{F}_Q \cdot \bar{F}_Q | B \rangle = -\frac{1}{2} \left( 1 + \frac{1}{N_c} \right)$$

(10)

$$\langle M | \bar{F}_Q \cdot \bar{F}_Q | M \rangle = -\frac{N_c^2 - 1}{2N_c}$$

(11)
At very short distances the potential between a heavy quark and antiquark should be described with perturbative QCD, and approximately given by an $N_c$-independent expression at leading order (LO) in terms of the running strong coupling $\alpha_s^{N_c}(r)$,

$$V_{Q\bar{Q}}^{N_c=3}\left(r\right) = -\frac{N_c^2-1}{8N_c^2} \frac{\alpha_s^{N_c}(r)}{r} = \frac{1}{12r} \frac{6}{r} \left(\frac{1}{11}\log(r\Lambda_{QCD})\right).$$  \tag{12}

At long distances it is of linear confining form and the corresponding string tension $\sigma$ is determined in LQCD. For $N_c = 3$ the Q potential has been computed in LQCD in the quenched approximation \cite{24}, and for $N_c > 3$ also \cite{7,8}. For $N_c = 3$, it is well described by the bosonic string model \cite{25}, namely:

$$V_{Q\bar{Q}}^{N_c=3}(r) = -\frac{\pi}{12r} + \sigma r. \tag{13}$$

The Coulomb term on the RHS is what results from the fluctuations of the string. It is remarkable that it provides the bulk of the Coulomb interaction down to the lattice spacings used in present day calculations. Using $\Lambda_{QCD}/\sqrt{\sigma} = 0.503(2)(40) + 0.33(3)(3)/N_c^2 + O(N_c^{-4})$ obtained in \cite{26} one gets that at $r\sqrt{\sigma} \sim 0.2$ the 1/r term in Eqs. (12) and (13) coincide. For the heavy quark mass corresponding to Compton wavelength much smaller than present lattice spacings, where the long distance potential plays a minor role, the Coulomb interaction will increasingly become the one predicted by perturbative QCD, Eq. (12).

At arbitrary $N_c$, $V_{Q\bar{Q}}^{N_c}$ will only receive corrections $O(1/N_c^2)$, as required by the $1/N_c$ expansion in pure gluodynamics. Assuming the leading scaling in $N_c$ for $\alpha_s$ and $\sigma$, and Eq. (11), the potential becomes:

$$V_{Q\bar{Q}}^{N_c}(r) = \frac{9}{8} \frac{N_c^2-1}{N_c^2} \frac{\alpha_s}{r} = \left(1 + O\left(1/N_c^2\right)\right) V_{Q\bar{Q}}^{N_c=3}(r). \tag{14}$$

This $N_c$ dependence will be loosely named “Casimir scaling”. This is verified by the ‘t Hooft coupling $\lambda = 4\pi N_c\alpha_s$ used in Refs. \cite{7,8}. Clearly this follows only if the above assumption is made, and with the present calculation at $N_c > 3$ it can be verified, as discussed below.

As mentioned earlier, the $1/N_c$ expansion requires definition because it compares different theories. The most obvious way to proceed is to require that certain quantities are independent of $N_c$, e.g., the string tension and quark masses at a given scale. Since the LQCD results of Refs. \cite{7,8} have the property that the string tension is approximately independent of $N_c$, i.e., $\sigma = -\frac{3}{8} N_c^{-1} \sigma_3 \sim$ const, this condition is adopted in what follows. The result from Fig. 1 vividly shows the $N_c$ independence of the $Q\bar{Q}$ potential within the current lattice uncertainties and the astonishing agreement with the bosonic string model \cite{25}. Thus, generalizing the $N_c$ lattice findings \cite{24} for all $N_c$ will be taken to be:

$$V_{Q\bar{Q}}^{N_c}(r) = V_{Q\bar{Q}}^{N_c=3}(r) = \frac{\pi}{12r} + \sigma r \tag{15}$$

From Eqs. (10)-(15) the two-body interaction potential in the baryon becomes:

$$V_{Q_1\bar{Q}_2\bar{Q}_3}^{N_c}(r) = \frac{V_{Q_1\bar{Q}_2}(r)}{N_c-1} = \frac{1}{N_c-1} \left(\frac{\pi}{12r} + \sigma r\right) \tag{16}$$

3. Mean field approximation and beyond

3.1. Mean field approximation

The calculation for different values of $N_c = 3, 5, 7, \ldots$ of the baryon mass with the Hamiltonian Eq. (5) requires solving separate few body problems with their inherent technical complications. In the large $N_c$ limit, however, an important simplification arises as a mean field approach becomes valid. Due to its color antisymmetry, the $N_c$-quark wave function in the baryon must be totally symmetric under simultaneous permutations of position and spin-flavor indices. In the ground state of the baryon, the $N_c$-quark spatial wave function is given in the Hartree approximation by \cite{2}:

$$\psi(x_1, \ldots, x_N) = \prod_{i=1}^N \phi(x_i). \tag{17}$$

For a single baryon, the baryon mass $M_B = \langle \psi|H|\psi \rangle \equiv \langle H \rangle_\psi$ is given by:

$$M_B = N_c m_Q + N_c \int d^3x \frac{1}{2m_Q} |\nabla \phi(x)|^2 + \frac{N_c(N_c-1)}{2} \int d^3x d^3y |\phi(x)|^2 |\phi(y)|^2 V_{Q\bar{Q}}(x-y). \tag{18}$$

The large $N_c$ scaling becomes obvious after the relation, Eq. (14), is used. It is useful to define the effective mean field potential $\bar{V}(x)$ generated by $N_c-1$ quarks

$$\bar{V}(x) = (N_c-1) \int d^3y V_{Q\bar{Q}}(x-y) |\phi(y)|^2 = \int d^3y V_{Q\bar{Q}}(x-y) |\phi(y)|^2, \tag{19}$$

where the Casimir scaling assumption provided by Eq. (16) has been used. The mean field potential is the self-energy of a quark within the hadron which sees the remaining $N_c-1$ quarks (which are coupled into the anti-fundamental representation $\bar{F}$).

The mean field equations are then obtained by minimizing with respect to a normalized $\phi(x)$ leading to the eigenvalue problem:

$$-\frac{1}{2m_Q} \nabla^2 \phi(x) + \bar{V}(x) \phi(x) = \epsilon \phi(x). \tag{20}$$

3.2. Numerical and variational solution
\[\phi(r) = \left( \frac{2}{\pi b^2} \right)^{\frac{3}{2}} e^{-r^2/b^2}\]  
\[(21)\]
yields a good approximation to this solution and allows for a simple analytical discussion.\(^1\)

3.3. CM corrections and mass formula

One standard and well documented approach to the mean field approximation in nuclear physics is the violation of Galilean invariance \([15,14]\) which is a symmetry of the starting Hamiltonian, Eq. (5), namely the invariance under the boost operation with velocity \(v\), \(\psi(x_1,\ldots,x_N) \rightarrow e^{imv \times \sum_i \kappa_i} \psi(x_1,\ldots,x_N)\), which implies the energy of the moving system to be given by \(E(P) = M + P^2/2m_Q N_c\), where the rest mass differs from the inertial mass \(M \neq N_c m_Q\).

Since the interest here is to include \(1/N_c\) corrections in the calculation, it is important to build a wave function that is an eigenfunction of the momentum. This is achieved by implementing, e.g., the Peierls–Yoccoz projection method \([15,14]\). However, for the simple Gaussian single particle wave function, Eq. (21), this corresponds just to replace \(N_c \rightarrow N_c - 1\) in the kinetic energy contribution.

Thus, the projection becomes trivial to deal with, and one obtains for a moving baryon of momentum \(P\):

\[\hat{M}_B = N_c m_Q + \frac{P^2}{2m_Q} + \frac{3(N_c - 1)}{2b^2 m_Q} + \frac{N_c}{b\sqrt{\pi}} (-\lambda^2 + b^2 \sigma),\]
\[(22)\]
where \(\lambda^2 = \pi/12\). Minimizing with respect to \(b\) \((b_0)\) yields the baryon mass at rest. At large \(m_Q\), \(b_0\) and the baryon mass become:

\[b_0 = \frac{3\sqrt{\pi}}{\lambda^2 m_Q} (N_c - 1) \left( 1 - 9\pi \left( \frac{N_c - 1}{N_c} \right) \right)^{\frac{2}{3}} \sigma / \lambda^6 m_Q^2 + O\left(1/m_Q^2\right)\]

\[\hat{M}_B = N_c m_Q + \frac{P^2}{2m_Q} - \left(\frac{N_c}{N_c - 1}\right)^2 \left( \frac{\lambda^4 m_Q}{6\pi} - \frac{3\sigma}{\lambda^2 m_Q^2} \right) + O\left(1/m_Q^2\right),\]
\[(23)\]
which shows a delayed onset of the heavy quark regime due to large numerical factors, namely in \(b_0\) the factor \(9\pi \left(\frac{N_c - 1}{N_c}\right)^\frac{2}{3} \frac{\sigma}{m_Q^3} = 1575 \left(\frac{N_c - 1}{N_c}\right)^\frac{2}{3} \frac{\sigma}{m_Q^3}\) and in \(\hat{M}_B\) the factor \(\frac{3\sigma}{\lambda^2 m_Q^2} = 11.46 \frac{\sigma}{m_Q}\). Thus, one should expect relativity to play a role even for moderately heavy quarks.

3.4. Relativistic corrections

Of course, a full relativistic treatment implies particle creation as implied by locality, and Poincaré invariant Hamiltonian methods with a fixed number of particles exhibit well known features

\(^1\) In the \(\sigma = 0\) case one has \(\hat{M}_B - 3m_Q = -0.00034\alpha^2 m_Q\) \([27]\) vs \(\hat{M}_B - 3m_Q = -0.00031\alpha^2 m_Q\) from Eq. (21). For the case \(\sigma \neq 0\) more sophisticated ansätze were tried embodying better short and long distance behaviors, but improvement is at the per cent level since the quarks are located in the mid-range region. Discussion of several possibilities will be given elsewhere.

\(^2\) Semiclassical quantization methods provide an alternative after due attention to zero modes is paid \([15,14]\).
quarks one expects that in the baryon only n-body interactions with \( n \leq N_c \) are of any significance. For \( N_c = 3 \) there is a long history of studying the 3-quark interactions, where there are two competing alternatives to confining forces of quarks in baryons, the \( \Delta \) (pairwise triangle shape) and the \( Y \) (juncture shape) inspired by string models [31].

Three body interactions have been addressed perturbatively [32] for arbitrary \( N_c \). In the present case, the non-perturbative effect of 3-body interactions can be visualized with one example. Consider a 3-body potential of the form:

\[
V_3(x_1, x_2, x_3) = \sum_{i=1}^{3} v_3(x_i - X) \, d_{abc} \, \lambda^a \otimes \lambda^b \otimes \lambda^c, 
\]

where \( X \) is the CM position of the three quarks. The expectation value of \( V_3 \) in the baryon ground state at rest can be evaluated explicitly choosing \( v_3(r) = \frac{1}{N_c^2} \left( -\frac{\lambda^2}{b} + \sigma_3 b \right) \) where \( \lambda_3 \) and \( \sigma_3 \) are \( \mathcal{O}(N_c^0) \), one obtains for the Gaussian wave function:

\[
\langle V_3 \rangle = 2 \sqrt{\frac{3}{\pi}} \left( N_c - \frac{5}{N_c} + \frac{4}{N_c^2} \right) \left( -\frac{\lambda^2}{b} + \sigma_3 b \right),
\]

where the color matrix element for the baryon was used,

\[
d_{abc} \lambda^a \otimes \lambda^b \otimes \lambda^c = \frac{(N_c - 3)!}{N_c!} \left( N_c^3 - 5N_c + \frac{4}{N_c} \right).
\]

Note that the expectation value of the 2-body interaction Eq. (22) and the one of the 3-body interaction studied here have the same form except that their \( N_c \) scalings differ by terms which are of relative order \( 1/N_c^2 \). Therefore, the 3-body forces cannot be distinguished from the 2-body ones unless those higher order terms in the expansion are taken into account. This is in a sense direct consequence of the mean field approximation, which naturally “hides” the n-body nature of the interactions. Other n-body forces are in principle possible for a large \( N_c \) baryon, whose color structure is given by \( 1/N_c^{n-1} \, d_{a_1 \cdots a_n} \lambda^{a_1} \otimes \cdots \lambda^{a_n} \), where \( d_{a_1 \cdots a_n} \) is the rank \( n \) invariant symmetric tensor of \( SU(N_c) \). A simple calculation shows that they contribute to the baryon mass with an overall factor \( N_c/n! \), which implies that even for very large \( N_c \), n-body forces with \( n > 5 \) become very suppressed.

3.7. Hyperfine effects

The simple OGE potential contains hyperfine components \( \mathcal{O}(m_Q^2) \), which have implications on meson spectra (see e.g. Ref. [33]), as they contribute at \( \mathcal{O}(N_c^0) \) in mesons, but contribute to hyperfine splitting in baryons only at \( \mathcal{O}(1/N_c) \). They can be easily evaluated as perturbations using the wave function obtained here. A quick calculation generalizing the \( N_c = 3 \) result [34] to arbitrary \( N_c \) gives for the hyperfine mass shift:

\[
\delta M_{BB}^{HF}(S) = \frac{8}{3\sqrt{\pi}} \frac{\alpha_s^N(m_Q^2)}{m_Q^2 b^3} (S(S + 1) - \frac{3}{4} N_c).
\]

They play no role for the spin-weighted average baryon mass Eq. (2).

4. Towards relating to LQCD results

Following the motivation of this work, the aim here is to compare the mean field description including relativistic and CM corrections with results from LQCD. At present, the only available LQCD results for ground state baryon masses at several \( N_c \) values are those of Refs. [7,8] (slightly updated in Ref. [18]), where quenched calculations have been undertaken at several values of the quark mass and at \( N_c = 3, 5, 7 \). While the purpose there was to pursue the light quark limit, here the opposite situation is emphasized where simplifications are expected and the quenched approximation is better fulfilled.

As discussed earlier, the explicit \( N_c \) dependence is inferred from taking \( \sigma \) to be \( N_c \) independent. The lattice results displayed in Refs. [7,8,18] are given in lattice units, with a the lattice spacing. Using the form of the quark–quark potential the Sommer parameter \( r_1 \) is determined by the standard definition

\[
r_1^2 V' \, \langle r_1 \rangle = -1,
\]

yielding in the present case

\[
r_1^2 \sigma = 1 - \frac{\pi}{12}.
\]

This value, namely \( r_1 \sqrt{\sigma} = 0.859 \), is roughly valid for the LQCD calculations with \( N_c = 3, 5 \) and \( 7 \), where the respective results from Table I of Ref. [7] are 0.856(5), 0.850(4) and 0.845(2). Using the values of \( r_1/a \) in the same Table one obtains respectively \( \sqrt{\sigma} a = 0.219(2), 0.225(2) \), and 0.216(1). For the level of precision of the present comparison it is therefore sufficient to take \( \sqrt{\sigma} a = 0.22 \) for all \( N_c \). While the main goal of [7,8] was to pursue the lowest quark mass limit, some moderately high quark masses were included. These are now used to compare with the results of this work.

The numerical results are presented in Fig. 2. As expected, the relativistic effects start becoming significant for \( 2\sqrt{\sigma} \sim m_Q \). The lattice data of Refs. [7,8] stop at twice larger values, so it would be highly interesting to extend the lattice calculations to the non-relativistic regime, where the comparison of the approach used and LQCD becomes more realistic. Qualitatively, it seems that there is a trend of the model and the LQCD results towards some agreement. The LQCD qualitative feature that \( M_B(N_c)/m_Q \) increases with \( N_c \) is also shown by the model, although it is not in good quantitative agreement.

At this point it is important to mention the issues involved in comparing the model with LQCD results. The main obstacle
pertains to the quark mass $m_Q$. While in the model it is a parameter (identifiable with a constituent quark mass), in LQCD it is a genuine QCD parameter which depends on the renormalization scheme \[35,36\]. Thus, the comparison made in Fig. 2 is only qualitative as it assumes that the masses are identical. For a more rigorous comparison one could proceed in different ways. Perhaps the best one would be to consider the spectrum of baryons for each $N_c$, adjusting the relation between the model and LQCD quark mass to best fit the LQCD results. While this would give a very realistic comparison, it seems unlikely that the spectrum of quark baryons at $N_c < 3$ could be calculated in LQCD in the foreseeable future. An intermediate approach is to compare the differences $(M_B(N_c) - M_B(N_c') / N_c)$, also adjusting the masses to a best fit (as mentioned earlier, it can be done with the present available LQCD results, but due to the rather small $m_Q$ values of those results it still unrealistic to compare). A rigorous approach along the lines discussed here would entail the use of an effective non-relativistic QCD theory for the heavy baryon \[37\], similar to the one for heavy quarkonium \[38\], where in principle it is possible to relate within QCD the quark mass of the effective theory to that of LQCD. Finally, an approach with immediate physical meaning would be to write the heavy baryon masses as a function of the corresponding quarkonium masses avoiding in this way having to relate quark masses and the use of the string tension as a fundamental parameter. Unfortunately, there are no direct LQCD evaluations of quarkonium masses; their masses are calculated via the use of the LQCD determined potential, similarly to what has been done in the present work for baryons.

The mean field approximation is visualized through the mean field potential $\bar{V}(r)$ created by the $N_c - 1$ quarks, see Eq. (19). In the present case, for zero momentum states and the Gaussian profile, Eq. (21) one obtains:

$$
\bar{V}(r) = \frac{b_0}{\sqrt{2\pi}} e^{-\frac{r^2}{2b_0^2}} \left( \lambda^2 - \sigma \left( \frac{b_0^2}{4} + r^2 \right) \right) \frac{1}{r} \text{erf} \left( \frac{\sqrt{2}r}{b_0} \right),
$$

which is shown for illustration, in Fig. 3 for different values of $N_c$ and $m_Q$. Improvements to this behavior correct for long distance behavior and will be discussed in a forthcoming publication.

5. Conclusions

In the present work, a scheme is put forward where the large $N_c$ expansion of baryon masses in the lattice can be described in terms of the mean field approximation as originally advocated by Witten and $1/N_c$ corrections thereof. The quark–quark potential is assumed to follow Casimir scaling at arbitrary $N_c$ and hence proportional to the quark–antiquark potential, which to good accuracy as per current LQCD calculations is $N_c$-independent. This provides a universal $N_c$ independent scheme where the ratio of the baryon mass to $\sqrt{\lambda}N_c$ can be numerically evaluated.

It was shown that the corrections to the mean field energy are generically $O(\sqrt{\lambda}N_c)$, but become $O(N_c^1)$, when the mean field energy takes its minimum value. This accuracy is the result of the density of quarks in the baryon growing as proportional with $N_c$. Among the estimated corrections are the leading in $N_c$ relativistic $O(m_Q^2)$ and subleading $O(N_c^0)$ CM corrections. Hyperfine splittings are removed by suitably averaging over spin states. When compared with available LQCD calculations, the present results account within 20% for the dimensionless ratio $(M_B - N_c m_Q) / \langle N_c \sqrt{\lambda} \rangle$ which is of natural size. This is encouraging, as it suggests to push the LQCD calculations to heavier quark masses and also to refine the calculations in the present work. The comparison undertaken here is so far just qualitative, as discussed in Section 4, due to the ambiguities in matching the model calculations to LQCD. More progress in this regard is needed in order to draw more rigorous comparisons.

One of the obvious benefits of the present investigation is the possibility of going beyond the ground state and extend these ideas to the excited baryon spectrum, where lattice calculations are admitted more involved and less accurate. LQCD calculations of excited baryons for $N_c > 3$ may still be an unreachable goal. However, it is likely that this will be achieved first with heavy quarks, and in that case the approach followed here can be easily used to make predictions of excited states. Finally, other heavy baryon properties, such as form factors, are easily derived with the wave functions obtained here.

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