ELSEVIER
SCIENCE DIRECT
PHYSICS LETTERS B
Physics Letters B 581 (2004) 175–181
www.elsevier.com/locate/physletb

Chiral soliton models, large N_c consistency and the Θ^+ exotic baryon

Thomas D. Cohen

Department of Physics, University of Maryland, College Park, MD 20742-4111, USA

Received 19 September 2003; received in revised form 26 November 2003; accepted 26 November 2003

Editor: W. Haxton

Abstract

Predictions for a light collective Θ^+ baryon state (with strangeness +1) based on the collective quantization of chiral soliton models are shown to be inconsistent with large N_c QCD. The lightest strangeness +1 state to emerge from the analysis has an excitation energy which at large N_c scales as N_c^0 while collective quantization is legitimate only for excitations which go to zero as $N_c \to \infty$. This inconsistency strongly suggests that predictions for Θ^+ properties based on collective quantization of chiral solitons are not valid.

There has been considerable recent excitement in hadronic physics. Several experimental groups have announced the identification of a narrow baryon resonance with a strangeness of +1 (i.e., containing one excess strange antiquark) [1]. Such a state is manifestly exotic in the sense of the quark model—it cannot be a simple three-quark state. This discovery has prompted considerable theoretical interest. Much of the theory has been in the context of generalized quark models in which the new baryon is identified as a pentaquark [2–10]. Unfortunately, the nature of this analysis is highly model dependent—there is no obvious way to see how phenomenological quark models emerge from QCD—and thus probably should be regarded presently as somewhat speculative. One theoretical approach to the problem clearly stands out—the analysis based on the SU(3) chiral soliton model treated with collective quantization [11–15]. This analysis has three obvious virtues: (i) the calculation predates the observation [11,12]; (ii) it made a strikingly accurate prediction of the mass [11,12] and has predicted a narrow width [12] consistent with those presently observed [16]; and (iii) although apparently based on a particular model—the chiral soliton model—the analysis is completely insensitive to the details of the model such as the profile function which emerges from the detailed dynamics.

This third point is particularly important. There has been considerable experience over the years with relations in chiral soliton models which are independent of the dynamical details going back nearly twenty years [17]. Typically such relations are exactly satisfied in the large N_c limit of QCD; the relations are derivable directly from large N_c consistency relations [18–21]. This holds for relations of typical static ob-

E-mail address: cohen@physics.umd.edu (T.D. Cohen).

0370-2693 © 2003 Elsevier B.V. Open access under CC BY license.
doi:10.1016/j.physletb.2003.11.065

servables (such as magnetic moments or axial couplings) considered in Ref. [17] and also for more esoteric quantities such as the nonanalytic quark mass dependence of observables near the chiral limit [22,23] or meson–baryon scattering observables [24]. Thus, it seems plausible that the analysis of Refs. [11–15] is similarly model-independent.

At first blush, this is quite satisfying: it appears that the observed Θ^+ state can be easily understood in terms of large N_c QCD and SU(3) flavor. The issue addressed in this Letter is whether this is, in fact, true. Despite the remarkable phenomenological success in predicting the mass and width of the Θ^+ seen in Ref. [12], a priori there is a compelling reason to doubt the validity of the analysis. Surprisingly this reason is *not* that the predicted state is a large N_c artifact but is associated with a more basic issue with soliton quantization. Here it is found that the prediction for the Θ^+ arises due to an inconsistent implementation of large N_c scaling in the soliton model; the prediction is an artifact of the treatment of collective variables in the model. In particular, it is shown here that the prediction depends on using collective quantization of the soliton outside the regime of validity of this method: states with positive strangeness such as the Θ^+ necessarily have an excitation energy of order N_c^0 while the semi-classical quantization method used to predict the state is only valid for excitations of order N_c^{-1}. An alternative argument based on general features of baryon states in large N_c QCD also indicates that the predicted Θ^+ state is spurious.

Let us begin by briefly reviewing the essential aspects of the analysis of Refs. [11–15]. The starting point is a treatment of SU(3) chiral soliton models which was developed in the mid-1980s [25]. In this approach one finds a classical static "hedgehog" configuration in an SU(2) subspace (the u–d subspace). The details of the profile are model dependent but the general structure of the theory is not. If one neglects SU(3) symmetry breaking effects then there are eight collective (rotational) variables which are then quantized semi-classically using an SU(3) generalization [25] of the usual SU(2) collective quantization scheme [26]. The collective Hamiltonian is given by

$$H^{\rm rot} = \frac{1}{2I_1}\sum_{A=1}^{3}\hat{J}_A'^2 + \frac{1}{2I_2}\sum_{A=4}^{7}\hat{J}_A'^2, \tag{1}$$

where I_1 (I_2) is the moment of inertia within (out of) the SU(2) subspace and $\hat{J}'_A$ are generators of SU(3) in a body-fixed (co-rotating) frame. Again, the numerical values of the moments of inertia are model dependent but the structure is not. There is an additional quantization constraint

$$J'_8 = -\frac{N_c B}{2\sqrt{3}}, \tag{2}$$

where B is the baryon number.

The explicit factor of N_c in Eq. (2) plays a central role in this Letter and it is useful to understand its origin. In Skyrme type models it follows directly from the Witten–Wess–Zumino term (which topology fixes to be an integer that can be identified with N_c). It can also be easily understood at the quark level. In a body-fixed frame the baryon number is associated with the SU(2) sub-manifold. There is also a body-fixed hypercharge associated with this sub-manifold which is related to the SU(3) generator in the usual manner: $Y' = -2J'_8/\sqrt{3}$. There is a general relation relating the baryon number, hypercharge and strangeness at large N_c which is valid at arbitrary N_c:

$$Y = \frac{N_c B}{3} + S, \tag{3}$$

this only coincides with the familiar relation $Y = B + S$ for $N_c = 3$. Eq. (3) follows from the fact that the hypercharge of up, down and strange quarks as being $1/3$, $1/3$ and $-2/3$, respectively. (These are the standard hypercharges of quarks in an $N_c = 3$ world. These hypercharge assignments must hold for general N_c provided hypercharge is isosinglet and traceless in SU(3) and has the property that the hypercharge of mesons is equal to the strangeness.) Given the fact that all three flavors of quark all have baryon number of $1/N_c$ while the strangeness is zero for u and d quarks and -1 for s quarks, one sees that Eq. (3) must hold. To complete the derivation of Eq. (2), note that in a body-fixed frame, the SU(2) sub-manifold has zero strangeness; accordingly Eq. (3) implies that $Y' = N_c B/3$ and the quantization condition in Eq. (2) immediately follows.

The masses which emerge from this depend on the quadratic Casimir of the SU(3) multiplet, $C_2 =$

$\sum_{A=1}^{8} \hat{J}_A^2$, and the angular momentum, J:

$$M_{\mathrm{SU}(3)} = M_0 + \frac{C_2}{2I_2} + \frac{(I_2 - I_1)J(J+1)}{2I_1 I_2} - \frac{N_c^2}{24 I_2},$$
$$\text{with } C_2 = \bigl(p^2 + q^2 + pq + 3(p+q)\bigr)/3, \tag{4}$$

where M_0 is a common soliton mass. C_2 is the quadratic Casimir and is expressed in terms of labels p, q which denote the SU(3) representation. The quantization condition in Eq. (2) greatly restricts the possible SU(3) representations: only SU(3) representations which contain hypercharge equal to $N_c/3$ are allowed: if the hypercharge in a body-fixed frame satisfies Eq. (2), the representation will include a state with that hypercharge. Moreover, since in the SU(2) manifold $I = J$ and $S = 0$, it follows that the number of angular momentum states associated with a representation, $2J + 1$, must equal the number of states in the representation with $S = 0$ (or equivalently with $Y = N_c/3$.

There is an ambiguity in how one implements this quantization. One might choose to quantize the theory at large N_c and then systematically put in $1/N_c$ corrections. Alternatively, in implementing the quantization condition of Eq. (2) one can fix $N_c = 3$ at the outset. To the extent that $N_c = 3$ can be considered large it ought not make any difference which of these approaches is used, provided that one is studying states which are not large N_c artifacts. Historically the choice of taking $N_c = 3$ at the outset has been standard [25]. Making this choice, it is straightforward to see that the lowest-lying states in this treatment are:

$$J = 1/2: \quad (p, q) = (1, 1) \quad (\text{octet}),$$
$$J = 3/2: \quad (p, q) = (3, 0) \quad (\text{decuplet}),$$
$$J = 1/2: \quad (p, q) = (0, 3) \quad (\text{anti-decuplet}). \tag{5}$$

The decuplet and the anti-decuplet can then be seen to have mass splittings relative to the octet given by

$$M_{10} - M_8 = \frac{3}{2I_1}, \tag{6}$$

$$M_{\overline{10}} - M_8 = \frac{3}{2I_2}. \tag{7}$$

The preceding analysis is a variant of quite standard 1980's vintage soliton physics. Note that this standard analysis of SU(3) solitons is only justified in the large N_c limit which plays an essential role in two ways. It justifies the use of the classical static hedgehog configurations; effects of quantum fluctuations around the hedgehogs are suppressed by $1/N_c$. It also justifies the semi-classical treatment in collective quantization; coupling between the collective motion and the internal structure of the hedgehog is also suppressed by $1/N_c$. It should be clear from the previous comment, however, that the validity of the collective approach depends on restricting its application to quantum collective modes. In order to track the N_c counting of various expressions we note that the moments of inertia $I_{1,2}$ scale as N_c.

The regime of validity of collective motion is critical to the analysis here, so it is useful to specify what it is and where it comes from. The key point is that a collective description is valid only for motion which is slow compared to the vibrational modes which are of order N_c^0. The vibrational modes are computed against a backdrop of a static soliton. This is valid providing the physical scale of the vibration is fast compared to the scale over which the soliton rotates. If this is not true one cannot separate the collective from the vibrational motion; in such a case the energy of the vibrational and collective motion are not additive and, indeed, it is a misnomer to refer to it as "collective" motion. Now the characteristic time scale of some type of quantized collective motion is given by the typical quantum mechanical result $\tau \sim (\Delta E)^{-1}$, where ΔE is the splitting between two neighboring collective levels. Thus collective motion is valid only for motion for which ΔE goes to zero in the large N_c limit.

Conventional treatments of collectively quantized SU(3) solitons identify the octet and decuplet states with the physical $N_c = 3$ octets and decuplets familiar from baryon spectroscopy, while the anti-decuplet has been dismissed as a large N_c artifact in much the same way that $I = J = 5/2$ baryons are generally dismissed as artifacts in SU(2) soliton models [26]. The principal intellectual argument of Ref. [12] is that the anti-decuplet should not be dismissed as a large N_c artifact. It argues that the anti-decuplet for SU(3) solitons can be distinguished from the $J = I = 5/2$ baryons in SU(2) in an essential way: the $J = I = 5/2$ baryon width would be predicted to be so wide with real world parameters that the state could not be observed [27]. In contrast, the anti-decuplet state might be expected to be narrow owing to suppressed phase space associated with the increased mass of kaons relative to pions. The fact that at the end of the calculation the predicted width of the Θ^+ is seen to be small is taken as a self-consistent justification of this approach.

Before proceeding further, a brief remark about the calculation in Ref. [12] is in order. Much of the detailed analysis concerns implementing SU(3) symmetry breaking effects in the calculation and how to fit the resulting parameters from data. For the present purposes, however, these are side issues. The central question of principle is whether the predicted collective anti-decuplet states are physical.

There is a very general argument why quantum number exotic collective states in chiral soliton models are expected to be spurious. A modern view of such models is that they encode the predictions of large N_c QCD relating the spin and flavor dependence of various observables [21]. The detailed numbers emerging from the models—the values of the masses, coupling constants and the like—are not reliable even at large N_c but the relations between them are. It is precisely because the analysis of Refs. [11–15] does not depend on dynamical details but merely on the structure of the collective quantization, that one might believe that it correctly encodes the underlying QCD physics. However, there is an alternative method to deduce the spin-flavor properties of large N_c baryons in a model independent way via the use of consistency conditions in describing meson–baryon scattering [21]. The results are well known: a contracted $\mathrm{SU}(2N_f)$ symmetry emerges in the large N_c limit. Baryon states fall into multiplets of $\mathrm{SU}(2N_f)$ and the low-lying states in these multiplets are split from the ground state by energies of order $1/N_c$—these excitations with the $\mathrm{SU}(2N_f)$ multiplets are collective. Moreover, the multiplet of low-lying baryons has been explicitly constructed—it coincides exactly with the low spin states of a quark model with N_c quarks confined to a single s-wave orbital [21]. Thus, it is well known that there are no low-lying collective baryon states in large N_c QCD with quantum numbers which are exotic for the large N_c world. In particular, there are no collective states with strangeness $+1$ in large N_c QCD. Any model which predicts such a collective state appears to be inconsistent with large N_c QCD.

This general argument strongly suggests that any strangeness $+1$ state predicted via collective quantization of a chiral soliton must somehow be spurious. Yet, at first glance, the derivation of Eq. (7) appears to be based on standard chiral soliton analysis. The issue is what, if anything, is wrong with the analysis? The answer lies in the collective quantization. Although the collective quantization of SU(3) solitons along the lines of [25] is the standard for the field, apparently, there has never been a careful study of the conditions for which the approach is consistent with large N_c QCD. As will be shown below, the approach appears to give excitations consistent with large N_c QCD for the lowest-lying $J = 3/2$ states but *not* for the exotic strangeness $+1$ states.

As stressed previously, the standard semi-classical treatment for collectively quantizing the solitons can only be justified in the large N_c approximation. The analysis outlined above appears to respect the underlying large N_c dynamics, at least formally. After all, the mass splitting in Eq. (7) goes as $1/I_2 \sim 1/N_c$. Thus, in the large N_c limit the splitting appears to become small which seems to imply that the motion is collective. The semi-classical quantization approach thereby looks to be justified self-consistently.

However, this is misleading: one can only see this collectivity clearly in the large N_c limit of the theory. Recall, however, that Eq. (7) was not derived in the large N_c limit. Its derivation depended on implementing the quantization condition in Eq. (2) with $N_c = 3$ at the outset. It was suggested above that making such a choice was innocuous, and indeed it is, *provided the states being studied are not artifacts*. However, since the entire question of relevance here is whether the states are spurious, we cannot start by using Eq. (7) to see if the motion is truly collective. Rather, one must study the full theory in its large N_c limit to see whether the motion turns out to be collective.

There are well-known peculiarities in studying SU(3) baryons in the large N_c limit. First and foremost among these is the fact that the SU(3) representations which emerge are not the ones we are familiar with at $N_c = 3$; indeed, as $N_c \to \infty$ all of these SU(3) representations become infinite-dimensional [21]. However, this presents no insurmountable problem phenomenologically, one simply associates those states in the representation with isospin and strangeness quantum numbers that survive down to the $N_c = 3$ with their real world analogs. The highly successful phe-

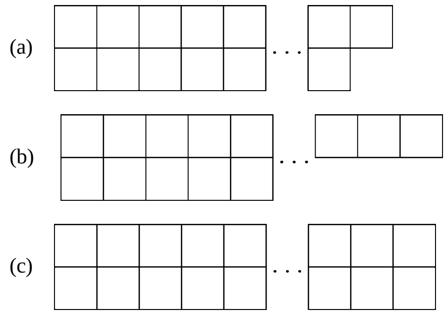

Fig. 1. Young tableau for arbitrary but large N_c: (a) the "8" representation with $(p,q) = (1, \frac{N_c-1}{2})$; (b) the "10" representation with $(p,q) = (3, \frac{N_c-3}{2})$; (c) the "$\overline{10}$" representation with $(p,q) = (0, \frac{N_c+3}{2})$. The Young tableau in (a) and (b) have N_c boxes; the tableau in (c) has $N_c + 3$ boxes.

nomenological study by Jenkins and Lebed of baryon masses based on large N_c scaling and SU(3) symmetry and its breaking was based precisely on this approach [28].

Consider the implementation of Eqs. (2) and (4) for N_c arbitrary and large. To ensure that our baryons remain fermions we restrict our attention to N_c odd. The lowest-lying representation compatible with Eq. (2) is easily seen to be $(p,q) = (1, \frac{N_c-1}{2})$ with $J = 1/2$ and is represented by the Young tableau (a) in Fig. 1. The states in this representation include those in the usual octet (and are thus taken to be their large N_c generalization); for convenience this representation will be denoted "8". The quotation marks serve to remind us that this is not really an octet. The next representation is $(p,q) = (3, \frac{N_c-3}{2})$ with $J = 3/2$; it is represented by the Young tableau (b) in Fig. 1 and is denoted by "10". Using Eq. (4), it is straightforward to see that:

$$M_{\text{“}10\text{”}} - M_{\text{“}8\text{”}} = \frac{3}{2I_1}. \tag{8}$$

Note that this is identical to the analogous result for the decuplet-octet splitting in Eq. (6). The significant point, however, is that since I_1 scales as N_c, this splitting *does* go to zero at large N_c indicating that the motion is, in fact, collective and thereby self-consistently justifying the use of collective quantization.

Next consider a large N_c representation analogous to the $\overline{10}$. The salient feature of the $\overline{10}$ representation is that it includes a state with strangeness $+1$. Thus, its large N_c analog should be taken to be the lowest-lying representation that includes a state with strangeness $+1$. This representation is $(p,q) = (0, \frac{N_c+3}{2})$ with $J = 1/2$; it is represented by the Young tableau (c) in Fig. 1 and is denoted as "$\overline{10}$". The excitation energy is given by

$$M_{\text{“}\overline{10}\text{”}} - M_{\text{“}8\text{”}} = \frac{3+N_c}{4I_2}. \tag{9}$$

Of course, Eq. (9) coincides with Eq. (7) for the special case of $N_c = 3$. However, unlike Eq. (7), Eq. (9) allows one to study the N_c scaling of the predicted splitting. Note that there is an explicit N_c in the numerator of the right-hand side while the denominator is proportional to I_2 which scales as N_c. Thus, the scaling at large N_c is given by

$$M_{\text{“}\overline{10}\text{”}} - M_{\text{“}8\text{”}} \sim N_c^0. \tag{10}$$

In the large N_c limit this splitting does not go to zero: the excitation is *not* collective. Note that the scaling in Eq. (10) is generic for states in large N_c QCD which are quantum number exotic in the sense that their quantum numbers cannot be obtained from N_c valance quarks. It is noteworthy that the *only* states whose excitation energies are of order N_c^{-1} are those whose Young tableau contains exactly N_c boxes; these are precisely the one seen in the general model independent analysis of Ref. [21].

Recall that the energy of the exotic Θ^+ was obtained using the collective quantization *which is only valid for collective modes*. However, as seen in Eq. (10), it is used to predict an excitation which is clearly not collective—its excitation energy remains finite at large N_c. Thus, the prediction of the low-lying Θ^+ state is based on using collective quantization outside its domain of validity.

Let us now revisit the argument in Ref. [12] based on the predicted hadronic widths that the predicted anti-decuplet state should not be regarded as spurious. Note this argument distinguished between the widths of the predicted anti-decuplet and the $J = 5/2$ states (which are generally regarded as large N_c artifacts). From the perspective of this Letter, it should be clear that these two states are entirely different beasts. The $J = 5/2$ states are collective modes whose properties one can safely predict in a large N_c world. The sole issue for the predicted $J = 5/2$ states is whether they survive in extrapolating back from large N_c to the real world at $N_c = 3$. In contrast, the strangeness $+1$ exotic states are not collective even in the large N_c limit; treating them using collective quantization will give rise to spuriously low energy modes. In short, the $J = 5/2$ state is spurious because its prediction depends on taking the large N_c limit too seriously, while the collective Θ^+ state is spurious because its prediction depends on not taking the large N_c limit seriously enough. Thus, although the reasons for which one regards the $J = 5/2$ state as spurious do not apply to the anti-decuplet, the anti-decuplet is spurious for entirely different reasons.

In summary, the predicted Θ^+ baryon in Refs. [11–15] was obtained using collective quantization in a regime where collective quantization does not apply. It was shown that quantum number exotic states in large N_c QCD have excitation energies which are of order N_c^0 and thus are not collective. Accordingly, the prediction of the Θ^+ as a collective excitation should be regarded as being invalid; the fact that the predicted mass was so near to the observed mass must be regarded as fortuitous.

Of course, none of the arguments presented here indicate that chiral soliton models are intrinsically incapable of describing exotic states or indeed of doing a reasonable phenomenological job in describing the Θ^+ baryon. However, if exotic states do exist in this class of models, they must be obtained by methods which are suitable to describe excitations of order N_c^0 rather than N_c^{-1}. Such methods do exist. For example one can use linear response theory to describe mesons scattering from baryons [29]. In principle, an exotic Θ^+ state could emerge in such a picture as a resonant state of a kaon and an ordinary baryon. However, there is no general argument that an exotic resonance would be generated for all such models and the excitation energy of such a state, if it exists, is completely model dependent. This does not imply that such an analysis is useless. One important aspect of large N_c QCD is that it *correlates* predictions. In particular, the existence of one light strangeness $+1$ resonant state implies the existence of other strangeness $+1$ resonant states which differ in energy from it by of order $1/N_c$. While the arguments presented in this Letter show why the order N_c^0 splitting between the ground state and the exotic are unreliable, the order $1/N_c$ splittings between exotic states are reliable. These predicted new states are explored in Ref. [30].

Acknowledgements

The author acknowledges Rich Lebed for several insightful remarks about this work. The support of the US Department of Energy for this research under grant DE-FG02-93ER-40762 is gratefully acknowledged.

References

[1] T. Nakano, et al., Phys. Rev. Lett. 91 (2003) 012002;
V.V. Barmin, et al., hep-ex/0304040;
S. Stepanyan, et al., hep-ex/0307018;
J. Barth, et al., hep-ex/0307083.
[2] F. Stancu, D.O. Riska, hep-ph/0307010.
[3] S. Capstick, P.R. Page, W. Roberts, hep-ph/0307019.
[4] B.G. Wybourne, hep-ph/0307170.
[5] A. Hosaka, hep-ph/0307232.
[6] R.L. Jaffe, F. Wilczek, hep-ph/0307341.
[7] C.E. Carlson, C.D. Carone, H.J. Kwee, V. Nazaryan, hep-ph/0307396.
[8] K. Cheung, hep-ph/0308176.
[9] L.Ya. Glozman, hep-ph/0308232.
[10] R.D. Matheus, F.S. Navarra, M. Nielsen, R. Rodrigues da Silva, S.H. Lee, hep-ph/0309001.
[11] M. Praszalowicz, in: M. Jezabek, M. Praszalowicz (Eds.), Workshop on Skyrmions and Anomalies, World Scientific, Singapore, 1987, p. 112. This work was not published in a journal at that time. An updated analysis based on this work may be found in hep-ph/0308114.
[12] D. Diakonov, V. Petrov, M. Polyakov, Z. Phys. A 359 (1997) 305.
[13] H. Walliser, V.B. Kopeliovich, hep-ph/0304058.
[14] D. Borisyuk, M. Faber, A. Kobushkin, hep-ph/0307370.
[15] H.-C. Kim, hep-ph/0308242.
[16] It has been argued, however, that the width's reported in the experiments may be broadened by experimental issues associated with resolution. Comparison with previous data seems to suggest that the actual width may be much narrower. This argument is detailed in S. Nussinov, hep-ph/0307357.
[17] G.S. Adkins, C.R. Nappi, Nucl. Phys. B 249 (1985) 507.
[18] J.-L. Gervais, B. Sakita, Phys. Rev. Lett. 52 (1984) 87;
J.-L. Gervais, B. Sakita, Phys. Rev. D 30 (1984) 1795.
[19] R.F. Dashen, A.V. Manohar, Phys. Lett. B 315 (1993) 425;
R.F. Dashen, A.V. Manohar, Phys. Lett. B 315 (1993) 438.
[20] E. Jenkins, Phys. Lett. B 315 (1993) 441.
[21] R.F. Dashen, E. Jenkins, A.V. Manohar, Phys. Rev. 49 (1994) 4713;
R.F. Dashen, E. Jenkins, A.V. Manohar, Phys. Rev. D 51 (1995) 3697.
[22] T.D. Cohen, W.B. Broniowski, Phys. Lett. B 292 (1992) 5.
[23] T.D. Cohen, Phys. Lett. B 359 (1995) 23;
T.D. Cohen, Rev. Mod. Phys. 68 (1995) 609.
[24] T.D. Cohen, R.F. Lebed, Phys. Rev. Lett. 91 (2002) 012001;

T.D. Cohen, R.F. Lebed, Phys. Rev. 67 (2002) 096008, hep-ph/0306102.
[25] E. Guadagnini, Nucl. Phys. B 236 (1984) 35;
P.O. Mazur, M.A. Nowak, M. Praszalowicz, Phys. Lett. B 147 (1984) 137;
A.V. Manohar, Nucl. Phys. B 248 (1984) 19;
M. Chemtob, Nucl. Phys. B 256 (1985) 600;
S. Jain, S.R. Wadia, Nucl. Phys. B 258 (1985) 713.
[26] G. Adkins, C. Nappi, E. Witten, Nucl. Phys. B 228 (1983) 522.
[27] T.D. Cohen, D.K. Griegel, Phys. Rev. D 43 (1991) 3089.
[28] E. Jenkins, R.F. Lebed, Phys. Rev. D 52 (1995) 282.
[29] M. Karliner, M.P. Mattis, Phys. Rev. Lett. 56 (1986) 428;
M. Karliner, M.P. Mattis, Phys. Rev. D 34 (1986) 1991.
[30] T.D. Cohen, R.F. Lebed, hep-ph/0309150.