I review recent lattice results on the spectrum and structure of baryons. Limitations due to the quenched approximation and un-physically heavy up and down quarks are discussed, and interfaces between first principles studies of QCD (or approximations to it) and “model building” are highlighted.

1. Introduction

With new exciting experiments ongoing and planned at several facilities, distributed over three continents, the past few years have witnessed the emergence of a new arena of research on the borderline between particle and nuclear physics: hadron physics. These developments have been complemented by several break-throughs in our theoretical understanding of the long distance, non-perturbative regime of hadronic matter, directly from QCD. Here I will focus on the most recent lattice results.

Progress in the field has been slow but steady. Several “revolutions” in the lattice technology, for instance non-perturbative renormalisation and improvement programmes of action and operators, were often hardly identifiable for the outside observer, a feature that we share with many other fields of research. In the case of strong QCD the failure of immediate satisfaction is due to the highly non-trivial phenomenology emerging from a Lagrangian as simple as that of QCD and yet so colourful and complex that even now it keeps experimentalists and theoreticians very busy.

A major but all too often neglected application of lattice results is to establish validity range and applicability of phenomenological models or effective field theories directly from QCD and to calculate low energy parameters that appear in these models or approximations. In quite a few sit-
uations, lattice predictions have matured to a degree that allows direct confrontation with experiment. Triggered by new experimental programmes and improved interaction with other theoreticians, an increasing number of lattice practitioners became interested in what has now become the domain of hadronic physics. This more focused approach resulted in an accelerated development of the field over the past two years with more to come in the future. I will go through the present state of lattice studies, discuss the two major sources of error, and review results on spectrum and structure of baryons, in particular $N\gamma \rightarrow \Delta$ transition form factors, charge distributions, three-body forces and the hyperon spin content.

2. The state of the lattice

Simulation results are obtained at finite lattice spacings and have to be extrapolated to the continuum limit. Moreover, the space-time box size should be big enough to comfortably accommodate the hadronic scales of interest. The dominant systematic uncertainties of present-day lattice results however tend to be related to unrealistically heavy $u$ and $d$ quark masses, which typically have to be chirally extrapolated over a long distance from $m_\pi^2 > 0.25 m_\rho^2$ to the physical limit $m_\pi^2 \approx 0.03 m_\rho^2$, often combined with the quenched approximation, i.e. neglecting the polarisation effects of sea quarks on the QCD vacuum.

The latter approximation, which easily reduces the computational effort by a factor of $10^3$, is justified in the limit of the number of colours $N_c \rightarrow \infty$, however, it is not a priori clear whether $1/3 \ll 1$. Phenomenology suggests that this assumption is not as outrageous as it might sound at first. This seems substantiated by simulations of pure $SU(N_c)$ gauge theories for $N_c = 2, \ldots, 6$ as well as by a comparison of quenched lattice predictions of the light hadron spectrum with experiment: ratios of light hadron masses on the lattice indeed have at last been found to be inconsistent with the observed spectrum. The differences are typically smaller than 10 %.

However, there are cases in which “quenching” clearly matters. Quantum mechanical perturbation theory tells us that an $O(c^2)$ correction to an energy eigenvalue corresponds to an $O(c)$ correction on wave functions which hence can easily be changed by as much as 30 %. This will affect quantities like decay constants and parton distributions. The situation with for instance bottomonium fine structure splittings or electroweak decay rates is even worse, as these are roughly proportional to the square of the wave function at the origin. Strong decays do not exist in this ap-
proximation as quarks and anti-quarks cannot annihilate, and therefore one might hesitate to expect very broad resonances to be modelled in a reasonable way. The same also holds true for any valence quark based model. We shall discuss some other shortcomings of quenching below.

Fortunately, with several new supercomputer installations on the horizon by early 2004 we are at the brink of the fourth generation of simulations including sea quarks. These certainly will bring us down in the quark mass to about the point where the $\rho$ meson becomes unstable and will at the same time allow for controlled continuum limit extrapolations.

Recently a way has been found to implement the analogue of exact chiral symmetry without “doublers” on the lattice, known as overlap, Neuberger or Ginsparg-Wilson (GW) fermions. One particular (truncated) realisation of the Neuberger operator are domain wall fermions. GW fermions are the only known consistent way of regularising a quantum field theory with massless fermions (like e.g. the standard model) and hence obviously a major theoretical break-through in particle physics as a whole. The famous Nielsen-Ninomiya no-go theorem is circumvented by relaxing the definition of chiral symmetry. I illustrate the yearly number of citations of these two important articles published in 1981 and 1982, respectively, in Figure 1.
and leave judgement to the reader which paper would have better impacted on Research Assessment Exercises. In practical terms GW fermions mean that realistically light quark masses will become accessible to future numerical simulations, first quenched\textsuperscript{10} and subsequently un-quenched.

Going to light quark masses is expensive as the linear spatial lattice extent $L$ has to be large relative to the inverse pion mass, to avoid strong finite size effects from pions exchanged “around” the boundaries. For instance if we intend to simulate a physical pion with say $L > 4/m_\pi$ this implies $L > 5.7$ fm. The main limitation however is that the effort of inverting the lattice Dirac operator scales with a large power of the inverse quark mass. This will improve with GW quarks that do not suffer from unphysical zero or near-zero modes: while at $m_\pi > 0.5 m_\rho$ GW fermions are easily two orders of magnitude more expensive to simulate than “standard” formulations, there will be a reward at small quark masses. In the quenched approximation this improvement factor has already turned out to be infinite, rendering a pion lighter than 200 MeV from impossible to possible.\textsuperscript{10} Chiral fermions also greatly simplify operator mixing and renormalisation.

3. Chiral extrapolations and “quenching”

Based on the assumptions that QCD bound states are mesons and baryons, that there is a mass gap and spontaneous chiral symmetry breaking at zero quark mass, an effective low energy chiral Lagrangian can be derived in the spirit of the Born-Oppenheimer approximation. This Lagrangian will, to leading order, describe interactions between the (fast moving) massless Nambu-Goldstone pions and other hadrons. In nature quarks and thus pions are not massless and the leading mass corrections are formally of order $m_\pi/\Lambda_{\chi SB}$ where $\Lambda_{\chi SB} \approx 4\pi F_\pi > 1$ GeV. Of course in an effective field theory the number of terms explodes at higher orders and predictive power is eventually lost, unless an early truncation is possible.$^{11}$

Third generation lattice simulations with sea quarks were limited to unrealistically heavy pions, heavier than about 400 MeV. Only recently masses as low as 180 MeV have become possible in the quenched approximation,$^{10}$ albeit not yet at sufficiently large volumes. To allow for a first principles connection with phenomenology it is mandatory to establish an overlap region between lattice simulations and (leading order) chiral perturbation theory. In general the size of this window will depend on the observable in question. GW fermions help a lot since only in this case we have a variant of chiral symmetry at finite lattice spacings while strictly speaking in other
formulations (without doublers) a sensitive comparison can only be performed after extrapolation of lattice results to the continuum limit. While even $400 \text{ MeV} \ll \Lambda_{\chi SB}$ (modulo the ambiguity of $\ll$ vs. $\ll$) such a pion is still doomed to “see” some of the internal structure of the proton, with an inverse charge radius of about 250 MeV: it is doubtful that the quark and gluon nature of QCD can completely be ignored with such a “hard” pion probe. With broadening baryonic wave functions such issues are likely to be more critical in QCD with light sea quarks than they already are in the quenched approximation.

Naïvely, hadron masses are a polynomial in the quark mass, $m_q \propto m^2_\pi$. However, pion loops give rise to non analytic (NA) functional dependence on the quark mass. For instance the nucleon mass is given by,

$$m_N(m_\pi) = m_N(0) + c_2 m^2_\pi + c_3 m^3_\pi + \cdots$$  \hspace{1cm} (1)

The coefficients of the NA terms are related to phenomenological low energy constants. For instance $c_3 = -3g^2_A/(32\pi F^2_\pi)$. In the quenched approximation, which neglects disconnected fermion lines, the leading NA (LNA) term is not proportional to $m^2_\pi$ anymore but to $m^2_\pi \ln m_\pi$, known as a chiral log. Only very recently convincing signs of such logarithmic behaviour have been found, albeit not yet with the expected coefficients.

Recently, corrections to standard LNA terms due to massive pions (re-
sulting in a different pion cloud) have been modelled, in particular by the
Adelaide group. Unlike the chiral Lagrangian itself, such models cannot
be derived in any limit directly from QCD, and the predictive power has to
be digested with some caution. Figure 2 reveals that we have not yet made
contact with chiral perturbation theory (dotted curve, $c_3$ fixed). Nonethe-
less the data are well fitted by Eq. (1), with $c_3$ as a free fit parameter
(dashed curve), and the extrapolated value has dangerously tiny statisti-
cal errors. The Adelaide cloudy bag model (CBM) suggests a differ-
tional form that screens the LNA term at pion masses close to a bag
parameter, $\Lambda$. With a fixed phenomenological value of $\Lambda$ the solid curv-
e is obtained while a three parameter fit yields the dashed-dotted curve. By
treating the difference as a systematic uncertainty such models allow us to
arrive at more realistic error estimates. Alternatives to the CBM include
NJL based as well as Bethe-Salpeter based ideas.

The quenched approximation is sick in several ways. I already men-
tioned that many unstable particles become stable. It is not always entirely
clear when this “bug” is a virtue and in which cases not. A more serious
flaw is that unitarity is violated since quarks are not included in a consis-
tent way. This defect, which becomes visible in particular at small quark
masses and in the scalar sector, also reflects onto and can be understood
in terms of the chiral expansion. For instance chiral logs originate from
this source. Another related problem is the absence of the axial anomaly
since quarks do not feed back onto the glue. This leads to the spontane-
ous breaking of a greater flavour symmetry group to start with and results in
$n_F^2 - 1$ pions: the $\eta'$ becomes heavy only at small sea quark
masses, with possible implications onto the phenomenology of excited
baryon states: the mass splitting between the nucleon ($N = N^{1/2+}$) and
the $N^* = N^{1/2-}$ for instance is a consequence of chiral symmetry breaking
and its size is related to transitions $N^* \to N + \pi \to N^*$. The number
of pions should naturally be expected to seriously affect such observables.
(Note that also in the quenched approximation there is a pion cloud since
valence quarks can travel forwards and backwards and forwards again in
time. Pion exchange is possible too, as valence quarks are swapped.)

To this end it has recently been suggested to model quenched LNA
terms and to subtract these from the corresponding lattice simulation re-
sults, replacing them by the expected un-quenched counterparts. Needless
to say that many assumptions are required, in particular relations between
low energy parameters appearing in quenched and un-quenched chiral La-
grangians and in the CBM. However, in the absence of lattice results with
light sea quarks, playing around with such models provides us with valuable rough ideas of possible sizes of quenching effects. On the other hand, once quality un-quenched data become available this will positively feed back onto such models, providing insights impossible to gain from experiment alone where quark masses and the number of active flavours are fixed.

Determinations of nucleon parton distributions seem to suffer from extremely large uncertainties in the chiral extrapolation; in particular this is so for the leading moments. In a recent study of the LHPC/SESAM Collaborations\textsuperscript{19} of quark distributions including sea quarks almost perfect agreement was found with quenched results, obtained both by them and by QCDSF,\textsuperscript{20} down to \( m_\pi \approx 600 \) MeV. However, a naïve polynomial extrapolation of \( \langle x \rangle_{u-d} \) was found to overestimate the phenomenological value by as much as a factor 5/3. Such a factor can indeed be accounted for by modelling pion mass effects on the LNA contribution in a sensible way,\textsuperscript{21} indicating that the quark mass is the dominant (and a huge) source of uncertainty.

4. Spectrum

In the past two years there has been a vibrant flow of lattice publications on the spectrum of excited nucleon states.\textsuperscript{22} All but one\textsuperscript{23} have been based on the quenched approximation so far. Only the LHPC/UKQCD/QCDSF Collaborations,\textsuperscript{24} using the Wilson-Clover action, attempted a continuum limit extrapolation. Other strategies were implementations of improved actions like the \( D_{234} \) action by Lee \textit{et al.}\textsuperscript{25} or the FLIC action by Melnitchouk \textit{et al.}\textsuperscript{26}. The Riken-BNL group employed domain wall fermions to improve the chiral properties.\textsuperscript{27}

One feature that all these articles share is a level inversion between the positive parity Roper \( N'(1440) \) resonance and the negative parity \( N^*(1535) \), relative to experiment. It was not clear whether this was a quenching effect, hence shared with many models, most of which are essentially \textit{quenched} too. But maybe the resonance observed in nature did something funny, to which the quark-only operator used on the lattice to create the Roper did not couple well. A pessimist might also argue that such resonances are quite unstable which could very well affect their position in ways that are not easily foreseeable, prior to doing the \textit{real thing}.

However, new results obtained with GW fermions by the Kentucky-Washington group\textsuperscript{10} at lighter pion masses, \( m_\pi > 180 \) MeV, than ever realised before indicate an exciting behaviour: as the pion becomes lighter
than about 550 MeV the slope of the Roper data bends downwards, heading straight in the direction of the experimental value (cf. Figure 3). While at the lightest quark mass \( Lm_\pi \approx 2.7 \) and finite size effects are likely, everything with \( m_\pi^2 > 0.1 \text{GeV}^2 \) is pretty safe. There is overlap with other lattice studies\(^\text{24,26,27}\) in the region \( m_\pi^2 > 0.25 \text{GeV}^2 \) and one might wonder why the onset of this behaviour has not been detected before. Inspection shows, however, that within the statistical errors an effect of the expected size would not have appeared significant in these cases.

Certainly, once the independence of the Kentucky results from Bayesian fitting priors and lattice volume has been firmly established, these are very exciting news that tell model builders how transitions between heavy and light quark regions can look like: a prime example for the cross fertilisation between lattice QCD, models and experiment in this thriving area of research. The Kentucky group has investigated many additional baryonic channels and it is astonishing how good quenched QCD seems to explain the positions of these strongly decaying resonances, once the quark masses become sufficiently light, another unsolved miracle of nature.

![Figure 3. Chiral extrapolation of the \( N, N' \) and \( N^* \) masses.\(^\text{10}\)](image-url)
5. Structure

Exciting results on the transition matrix element for electromagnetic $N$ to $\Delta$ transitions have been reported. The relevant operator is usually decomposed into three Sachs type form factors, accompanying magnetic dipole transitions, $G_{M1}(q^2)$ as well as electric and Coulomb/scalar quadrupole transitions, $G_{E2}(q^2)$ and $G_{C2}(q^2)$. The quadrupole transitions ratios $R_{EM} = G_{E2}/G_{M1}$ and $R_{SM} = G_{C2}/G_{M1}$ are related to the question of nucleon deformation. Recent experimental results, in particular from the CLAS and OOPS Collaborations indicate the ratios $R_{EM} \approx -0.02$ and $R_{SM} \approx -0.07$ around $q^2 \approx 0.53$, which is the lowest momentum transfer that could be realised in a recent lattice study with two flavours of sea quarks, not much lighter than the strange. These first quantitative lattice results seem to get the sign right and point towards values $R_{EM} = -0.03(1)$ and $R_{SM} = -0.03(2)$, consistent with experiment. Interestingly, the quenched reference result $R_{EM} = -0.009(8)$ is compatible with zero, indicating that a more realistic pion cloud might be essential for the effect. Obviously, lattice simulations are still miles away from computing, say, octupole form factors for transitions between orbitally excited nucleons.

In another lattice investigation the electric charge and matter distributions of $\pi$, $\rho$, $N$ and $\Delta$ have been studied at rest. One can easily calculate the charge radius from such a distribution and the Fourier transform is related to form factors at zero momentum transfer. While the (cut-off and scheme dependent) matter distributions of all investigated hadrons were found to be remarkably similar there are differences in the charge distributions. In the quenched approximation the pion is only approximately half as wide as the $\Delta$, with $N$ and $\rho$ almost of the same size as the $\Delta$. Even with a pion mass of almost 600 MeV un-quenching significantly broadened the latter three distributions, and in particular the $\Delta$, while the $\pi$ remains completely unaffected. The quenched proton charge radius comes out to be $r_p = 0.59(4)$ fm, after polynomial extrapolation to the physical limit. This has to be compared with the value $r_p \approx 0.81$ fm from dipole fits to experimentally determined electromagnetic form factors. The above direct lattice result is in agreement with $r_p \approx 0.6$ fm obtained indirectly from a dipole fit to electric and magnetic proton form factors calculated on the lattice by the QCDSF Collaboration. However, from recent JLAB Hall A results, showing an energy dependence of the ratio $G_e/G_m$, we know that the dipole ansatz is not the whole truth, which stresses the importance of
determining quantities like charge radii in more than one way. The same QCDSF work suggests \( r_p \approx 0.7 \text{ fm} \) for two flavours of sea quarks, closer to the phenomenological value.

For the first time deformations of the hadron wave functions have been investigated.\(^{32}\) Obviously, a spectroscopic quadrupole moment becomes only possible for an angular momentum of at least one, such that deviations from spherical symmetry are ruled out for the nucleon. However, the \( \rho \) was reported to be somewhat prolate with the \( \Delta \) being slightly oblate and the deformations appeared to increase when including sea quarks. In neither of the cases was the deformation of the \( \Delta \) statistically significant though.

It is not clear whether QCD forces are \textit{stringy} with a three body potential that depends on the length of the shortest path connecting the three quarks to one point (Mercedes or \( Y \) law) or if they merely can be understood in terms of sums of two body interactions (\( \Delta \) law).\(^{35}\) Leaving aside the question whether an instantaneous Born-Oppenheimer interaction picture is appropriate for a situation with light quarks, the nucleon wave function seems to be better described by a Schrödinger equation with a potential that depends on the “\( \Delta \)-distance” than by the \( Y \) configuration.\(^{32}\)

This finding is also supported by recent lattice calculations of \( SU(3) \) and \( SU(4) \) gauge theory as well as of the \( Z_3 \) Potts model.\(^{36}\) The results indicate that for interquark distances of physical relevance, i.e. \( r < 1 \text{ fm} \), the \( \Delta \)-picture yields a very accurate description while, long after the “string” of QCD with sea quarks is broken, the \( Y \) picture is eventually approached.

Recently, spin and transversity of the quark contributions to the \( \Lambda \) hyperon have been calculated by QCDSF.\(^{38}\) The fraction \( \Delta q_A \) of the spin carried by quarks of flavour \( q \) is given by the forward matrix element of the axial vector current, \( \langle \Lambda(p', s)| \bar{q}\gamma_\mu \gamma_5 q | \Lambda(p, s) \rangle = 2 (p' - p)_\mu \Delta q_A \), i.e. by the lowest moment of the helicity distribution. Assuming \( SU(3) \) flavour symmetry one can express the \( \Delta q_A \)s in terms of proton spin fractions \( \Delta q_p \). Converting the lattice results into the \( \overline{MS} \) scheme at \( \mu = 2 \text{ GeV} \) yields an axial charge \( g_A = \Delta u_p - \Delta d_p \) that falls short of the experimental value \( g_A \approx 1.26 \) by 10-20 \%, due to quenching and/or due to uncertainties in the chiral extrapolation, such that we cannot expect the result \( \Delta u_\Lambda = \Delta d_\Lambda = -0.02(4) \) to be consistent with the real world either.

The lattice data on proton and \( \Lambda \) spin fractions turn out to be in excellent agreement with \( SU(3) \) flavour symmetry, an observation that might be independent of sea quarks and certainly consistent with the success in understanding semileptonic decay rates of the \( \Lambda \). Combining experi-
mental data obtained for the proton helicity with the systematic error on $SU(3)_F$ symmetry suggested by this lattice study, yields the predictions $\Delta u_\Lambda = \Delta d_\Lambda = -0.17(3)$ and $\Delta s_\Lambda = 0.63(3)$, again in the $\overline{MS}$ scheme at $\mu = 2$ GeV.

6. Outlook

Experimental efforts are paralleled by lattice studies of baryon spectrum and structure. The main limitation that has to be overcome in the near future is the use of unrealistically heavy quarks in quenched simulations, as well as in simulations with sea quarks. Recent theoretical developments and advances in computer technology mean that optimism is justified.

Acknowledgments

This work is supported by PPARC grants PPA/A/S/2000/00271, PPA/G/O/1998/00559 and EU grant HPRN-CT-2000-00145.

References

1. See e.g., R. F. Lebed, these proceedings [arXiv:hep-ph/0301279].
2. B. Lucini, M. Teper and U. Wenger, Phys. Lett. B545, 197 (2002) [arXiv:hep-lat/0206029]; B. Lucini and M. Teper, JHEP 0106, 050 (2001) [arXiv:hep-lat/0103027].
3. S. Aoki et al., arXiv:hep-lat/0206009.
4. G. S. Bali and P. Boyle, Phys. Rev. D59, 114504 (1999) [arXiv:hep-lat/9809180].
5. H. Neuberger, Phys. Lett. B417, 141 (1998) [arXiv:hep-lat/9707022].
6. D. B. Kaplan, Phys. Lett. B288, 342 (1992) [arXiv:hep-lat/9206013].
7. H. B. Nielsen and M. Ninomiya, Nucl. Phys. B185, 20 (1981).
8. P. H. Ginsparg and K. G. Wilson, Phys. Rev. D25, 2649 (1982).
9. The idea for such a Figure has been “stolen” from P. Hasenfratz: http://www.theorie.physik.uni-wuppertal.de/EU2001/slides/ph/ph.0.phtml
10. F. X. Lee, S. J. Dong, T. Draper, I. Horvath, K. F. Liu, N. Mathur and J. B. Zhang, arXiv:hep-lat/0208070; F. X. Lee, these proceedings.
11. See e.g., J. Nieves, these proceedings.
12. D. B. Leinweber, A. W. Thomas, K. Tsushima and S. V. Wright, Nucl. Phys. Proc. Suppl. 83, 179 (2000) [arXiv:hep-lat/9909109]; R. D. Young, D. B. Leinweber, A. W. Thomas and S. V. Wright, Phys. Rev. D 66, 094507 (2002) [arXiv:hep-lat/0205017].
13. A. Ali Khan et al., Phys. Rev. D65, 054505 (2002) [arXiv:hep-lat/0105015].
14. C. R. Allton et al., Phys. Rev. D65, 054502 (2002) [arXiv:hep-lat/0107021].
15. A. N. Kvinikhidze, M. C. Birse and B. Blankleider, Phys. Rev. C66, 045203 (2002) [arXiv:hep-ph/0110060].
16. P. J. Bicudo, G. Krein and J. E. Ribeiro, *Phys. Rev.* **C64**, 025202 (2001) [arXiv:hep-ph/0105289].
17. G. S. Bali *et al.*, *Phys. Rev.* **D64**, 054502 (2001) [arXiv:hep-lat/0102002]; T. Struckmann *et al.*, *Phys. Rev.* **D63**, 074503 (2001) [arXiv:hep-lat/0010005].
18. A. W. Thomas, arXiv:hep-lat/0208023.
19. D. Dolgov *et al.*, *Phys. Rev.* **D66**, 034506 (2002) [arXiv:hep-lat/0201021].
20. S. Capitani, M. Göckeler, R. Horsley, D. Pleiter, P. Rakow, H. Stüben and G. Schierholz, *Nucl. Phys. Proc. Suppl.* **106**, 299 (2002) [arXiv:hep-lat/0111012].
21. W. Detmold, W. Melnitchouk, J. W. Negele, D. B. Renner and A. W. Thomas, *Phys. Rev. Lett.* **87**, 172001 (2001) [arXiv:hep-lat/0103006].
22. D. Richards, these proceedings.
23. C. M. Maynard and D. G. Richards, arXiv:hep-lat/0209165.
24. M. Göckeler, R. Horsley, D. Pleiter, P. E. Rakow, G. Schierholz, C. M. Maynard and D. G. Richards, *Phys. Lett.* **B532**, 63 (2002) [arXiv:hep-lat/0206022].
25. F. X. Lee, D. B. Leinweber, L. Zhou, J. M. Zanotti and S. Choe, *Nucl. Phys. Proc. Suppl.* **106**, 248 (2002) [arXiv:hep-lat/0110164].
26. W. Melnitchouk *et al.*, arXiv:hep-lat/0202022; W. Melnitchouk, J. N. Hedditch, D. B. Leinweber, A. G. Williams, J. M. Zanotti and J. B. Zhang, arXiv:hep-lat/0201042.
27. S. Sasaki, T. Blum and S. Ohta, *Phys. Rev.* **D65**, 074503 (2002) [arXiv:hep-lat/0102010].
28. D. Drechsel, these proceedings.
29. K. Joo *et al.*, *Phys. Rev. Lett.* **88**, 122001 (2002) [arXiv:hep-ex/0110007]; R. Arndt, W. Briscoe, I. Strakovsky and R. Workman, arXiv:nucl-th/0301068, these proceedings.
30. C. Mertz *et al.*, *Phys. Rev. Lett.* **86**, 2963 (2001) [arXiv:nucl-ex/9902012].
31. C. Alexandrou *et al.*, arXiv:hep-lat/0209074.
32. C. Alexandrou, P. de Forcrand and A. Tsapalis, *Phys. Rev.* **D66**, 094503 (2002) [arXiv:hep-lat/0206026].
33. M. Göckeler, R. Horsley, D. Pleiter, P. E. Rakow and G. Schierholz, arXiv:hep-ph/0108105.
34. M. K. Jones, *Nucl. Phys.* **A699** (2002) 124.
35. G. S. Bali, *Phys. Rept.* **343**, 1 (2001) [arXiv:hep-ph/0001312].
36. C. Alexandrou, P. de Forcrand and O. Jahn, arXiv:hep-lat/0209062; C. Alexandrou, P. De Forcrand and A. Tsapalis, *Phys. Rev.* **D65**, 054503 (2002) [arXiv:hep-lat/0107006].
37. G. S. Bali *et al.*, *Phys. Rev.* **D62**, 054503 (2000) [arXiv:hep-lat/0003012]; C. McNeile, C. Michael and P. Pennanen, *Phys. Rev.* **D65**, 094505 (2002) [arXiv:hep-lat/0201006].
38. M. Göckeler, R. Horsley, D. Pleiter, P. E. Rakow, S. Schaefer, A. Schäfer and G. Schierholz, *Phys. Lett.* **B545**, 112 (2002) [arXiv:hep-lat/0208017].