

Nonperturbative Determination of Quark Masses in Quenched Lattice QCD with the Kogut-Susskind Fermion Action

S. Aoki,¹ M. Fukugita,² S. Hashimoto,³ K-I. Ishikawa,⁴ N. Ishizuka,^{1,5} Y. Iwasaki,^{1,5} K. Kanaya,^{1,5} T. Kaneda,¹
S. Kaya,³ Y. Kuramashi,³ M. Okawa,³ T. Onogi,⁴ S. Tominaga,³ N. Tsutsui,⁴ A. Ukawa,^{1,5}
N. Yamada,⁴ and T. Yoshié^{1,5}

(JLQCD Collaboration)

¹*Institute of Physics, University of Tsukuba, Tsukuba, Ibaraki 305-8571, Japan*

²*Institute for Cosmic Ray Research, University of Tokyo, Tanashi, Tokyo 188-8502, Japan*

³*High Energy Accelerator Research Organization (KEK), Tsukuba, Ibaraki 305-0801, Japan*

⁴*Department of Physics, Hiroshima University, Higashi-Hiroshima, Hiroshima 739-8526, Japan*

⁵*Center for Computational Physics, University of Tsukuba, Tsukuba, Ibaraki 305-8577, Japan*

(Received 26 January 1999)

We report results of quark masses in quenched lattice QCD with the Kogut-Susskind fermion action, employing the regularization independent scheme of Martinelli *et al.* to nonperturbatively evaluate the renormalization factor relating the bare quark mass on the lattice to that in the continuum. Calculations are carried out at $\beta = 6.0, 6.2,$ and $6.4,$ from which we find $m_{ud}^{\overline{\text{MS}}}(2 \text{ GeV}) = 4.23(29) \text{ MeV}$ for the average up and down quark mass and, with the ϕ meson mass as input, $m_s^{\overline{\text{MS}}}(2 \text{ GeV}) = 129(12) \text{ MeV}$ for the strange mass in the continuum limit. These values are about 20% larger than those obtained with the one-loop perturbative renormalization factor. [S0031-9007(99)09226-1]

PACS numbers: 12.38.Gc, 12.15.Ff

The values of quark masses are fundamental parameters of the standard model which are not directly accessible through experimental measurements. Lattice QCD allows their determination through a calculation of the functional relation between quark masses and hadron masses. For this reason a number of lattice QCD calculations have been carried out to evaluate quark masses, employing the Wilson, clover, or Kogut-Susskind (KS) fermion action [1]. These actions are different lattice discretizations of the continuum fermion action, and an agreement of quark masses calculated with them provides an important check on the reliability of lattice results.

An important ingredient in these calculations is the renormalization factor relating the bare lattice quark mass to that in the continuum. While perturbation theory is often used to evaluate this factor, uncertainties due to higher order terms are quite significant in the range of the QCD coupling constant accessible in today's numerical simulations. A nonperturbative determination of the renormalization factor is therefore necessary for a reliable calculation of quark masses, and effort in this direction has recently been pursued for the Wilson and clover fermion actions [2–5].

The need for a nonperturbative determination of the renormalization factor is even more urgent for the KS action since the one-loop correction [6,7] is as large as 50% in present simulations. In this Letter we report a study to meet this need [8,9]: we calculate the renormalization factor of bilinear quark operators for the KS action nonperturbatively using the regularization independent (RI) scheme of Ref. [2] developed for the Wilson and clover

actions. The results for the scalar operator, combined with our previous calculation of bare quark masses [10], lead to a nonperturbative determination of the quark masses in the continuum limit.

In the RI scheme, the renormalization factor of a bilinear operator \mathcal{O} is obtained from the amputated Green function,

$$\Gamma_{\mathcal{O}}(p) = S(p)^{-1} \langle 0 | \phi(p) \mathcal{O} \bar{\phi}(p) | 0 \rangle S(p)^{-1}, \quad (1)$$

where the quark two-point function is defined by $S(p) = \langle 0 | \phi(p) \bar{\phi}(p) | 0 \rangle$. The quark field $\phi(p)$ with momentum p is related to the one-component KS field $\chi(x)$ by $\phi_A(p) = \sum_y \exp(-ip \cdot y) \chi(y + aA)$, where $y_\mu = 2an_\mu$, $p_\mu = 2\pi n_\mu / (aL)$ with $-L/4 \leq n_\mu < L/4$ and $A_\mu = 0, 1$.

Bilinear operators have a form

$$\mathcal{O} = \sum_{yABab} \bar{\phi}_A^a(y) (\overline{\gamma_S \otimes \xi_F})_{AB} U_{AB}^{ab}(y) \phi_B^b(y), \quad (2)$$

where $(\overline{\gamma_S \otimes \xi_F})$ refers to Dirac (γ_S) and KS flavor (ξ_F) structure [7], and the indices a and b refer to color. The factor $U_{AB}^{ab}(y)$ is the product of gauge link variables along a minimum path from $y + aA$ to $y + aB$.

The renormalization condition imposed on $\Gamma_{\mathcal{O}}(p)$ is given by

$$Z_{\mathcal{O}}^{\text{RI}}(p) \cdot Z_{\phi}(p) = \text{Tr}[P_{\mathcal{O}}^{\dagger} \Gamma_{\mathcal{O}}(p)], \quad (3)$$

where $P_{\mathcal{O}}^{\dagger} = (\overline{\gamma_S^{\dagger} \otimes \xi_F^{\dagger}})$ is the projector onto the tree-level amputated Green function. The wave function renormalization factor $Z_{\phi}(p)$ can be calculated by the

TABLE I. Run parameters.

β	$1/a(\text{GeV})$	am			No. conf.
6.0	1.855(38)	0.010	0.020	0.030	30
6.2	2.613(87)	0.008	0.015	0.023	30
6.4	3.425(72)	0.005	0.010	0.020	50

condition $Z_V(p) = 1$ for the conserved vector current corresponding to $(\gamma_\mu \otimes I)$. Since the RI scheme explicitly uses the quarks in external states, gauge fixing is necessary. We employ the Landau gauge throughout the present work.

The relation between the bare operator on the lattice and the renormalized operator in the continuum takes the form

$$\mathcal{O}_{\overline{\text{MS}}}(\mu) = U_{\overline{\text{MS}}}(\mu, p) Z_{\text{RI}}^{\overline{\text{MS}}}(p) / Z_{\text{O}}^{\text{RI}}(p) \mathcal{O}, \quad (4)$$

where $U_{\overline{\text{MS}}}(\mu, p)$ is the renormalization-group running factor in the continuum from momentum scale p to μ . We adopt the naive dimensional regularization (NDR) with the modified minimum subtraction scheme ($\overline{\text{MS}}$) in the continuum. The factor $Z_{\text{RI}}^{\overline{\text{MS}}}(p)$ provides matching from the RI scheme to the $\overline{\text{MS}}$ scheme. These two factors are calculated perturbatively in the continuum. For our calculation of the quark mass we apply the relation (4) in the scalar channel in the chiral limit, i.e., $1/Z_m^{\text{RI}} = Z_S^{\text{RI}}$.

Our calculations are carried out in quenched QCD. Gauge configurations are generated with the standard plaquette action at $\beta = 6.0, 6.2$, and 6.4 on a 32^4 lattice. For each β we choose three bare quark masses tabulated in Table I where the inverse lattice spacing $1/a$ is taken from our previous work [10].

We calculate Green functions for 15 momenta in the range $0.038553 \leq (ap)^2 \leq 1.9277$. Quark propagators are evaluated with a source in a momentum eigenstate. We find that the use of such a source results in very small statistical errors of $O(0.1\%)$ in Green functions.

The RI method completely avoids the use of lattice perturbation theory. We do not have to introduce any ambiguous scale, such as q^* [11], to improve on one-loop results. An important practical issue, however, is whether the renormalization factor can be extracted from a momentum range $\Lambda_{\text{QCD}} \ll p \ll O(1/a)$ keeping under control the higher order effects in continuum perturbation theory, nonperturbative hadronization effects, and the discretization error on the lattice. These effects appear as p dependence of the renormalization factor in (4), which should be absent if these effects are negligible.

In Fig. 1 we compare the scalar renormalization factor $Z_S^{\text{RI}}(p)$ with that for pseudoscalar $Z_P^{\text{RI}}(p)$ for three values of bare quark mass am at $\beta = 6.0$. From chiral symmetry of the KS fermion action, we naively expect a relation $Z_S^{\text{RI}}(p) = Z_P^{\text{RI}}(p)$ for all momenta p in the chiral limit. Clearly this does not hold with our result toward small momenta, where $Z_P^{\text{RI}}(p)$ rapidly increases as $m \rightarrow 0$, while $Z_S^{\text{RI}}(p)$ does not show such a trend.

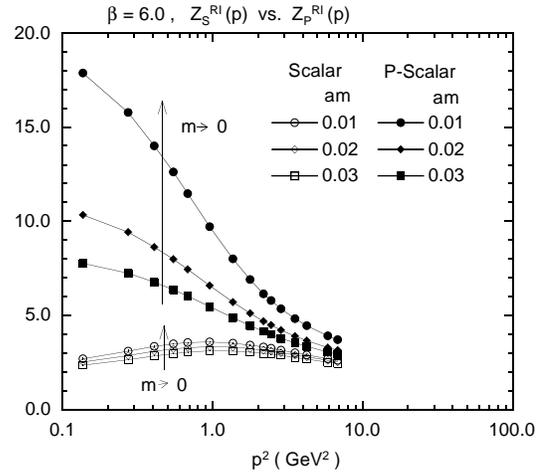


FIG. 1. The renormalization factor for the scalar $Z_S^{\text{RI}}(p)$ and the pseudoscalar $Z_P^{\text{RI}}(p)$.

To understand this result, we note that chiral symmetry of KS fermion action leads to the identities

$$Z_S^{\text{RI}}(p) Z_\phi(p) = \partial M(p) / \partial m, \quad (5)$$

$$Z_P^{\text{RI}}(p) Z_\phi(p) = M(p) / m, \quad (6)$$

with $M(p) = \text{Tr}[S(p)^{-1}]$. In Fig. 2 $M(p)$ in the chiral limit obtained by a linear extrapolation in m is plotted. It rapidly decreases for large momenta, but takes nonzero values for small momenta. Equation (6) implies that $Z_P^{\text{RI}}(p)$ diverges in the chiral limit for small momenta, which is consistent with the result in Fig. 1.

The function $M(p)$ is related to the chiral condensate as follows:

$$\langle \phi \bar{\phi} \rangle = \sum_p \text{Tr}[S(p)] = \sum_p \frac{M(p)}{\sum_\mu C_\mu(p)^2 + M(p)^2}, \quad (7)$$

where $C_\mu(p) = -i \text{Tr}[(\gamma_\mu \otimes I) S(p)^{-1}] / \cos(p_\mu a)$. A nonvanishing value of $M(p)$ for small momenta would lead to a nonzero value of the condensate. Therefore the divergence of $Z_P^{\text{RI}}(p)$ near the chiral limit is a

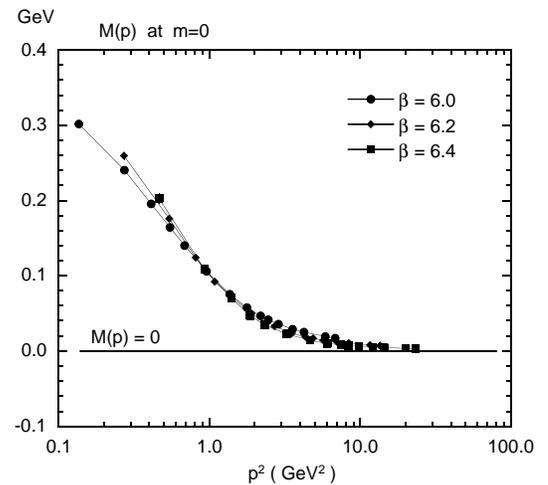


FIG. 2. $M(p)$ in the chiral limit.

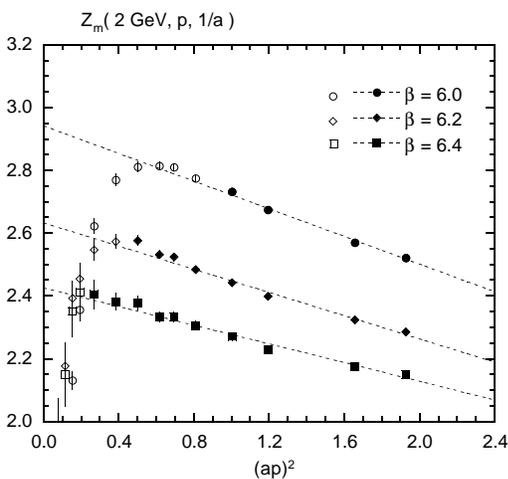


FIG. 3. The ratio $Z_m(\mu, p, 1/a)$ at $\mu = 2$ GeV. For each β the filled data points are used for linear interpolation in $(ap)^2$.

manifestation of spontaneous breakdown of chiral symmetry; it is a nonperturbative hadronization effect arising from the presence of massless Nambu-Goldstone boson in the pseudoscalar channel.

While we do not expect the pseudoscalar meson to affect the scalar renormalization factor $Z_S^{R1}(p)$, as indeed observed in the small quark mass dependence seen in Fig. 1, the above result raises a warning that $Z_S^{R1}(p)$ may still be contaminated by hadronization effects for small momenta.

In Fig. 3 we show the momentum dependence of $Z_m(\mu, p, 1/a) \equiv U_{\overline{MS}}(\mu, p)Z_{RI}^{\overline{MS}}(p)Z_S^{R1}(p)$ which is the renormalization factor from the bare quark mass on the lattice to the renormalized quark mass at scale μ in the continuum. Here we set $\mu = 2$ GeV and use the three-loop formula [12] for $U_{\overline{MS}}$ and $Z_{RI}^{\overline{MS}}$. While $Z_m(\mu, p, 1/a)$ should be independent of the quark momentum p , our results show a sizable momentum dependence which is almost linear in $(ap)^2$ for large momenta (filled symbols in Fig. 3).

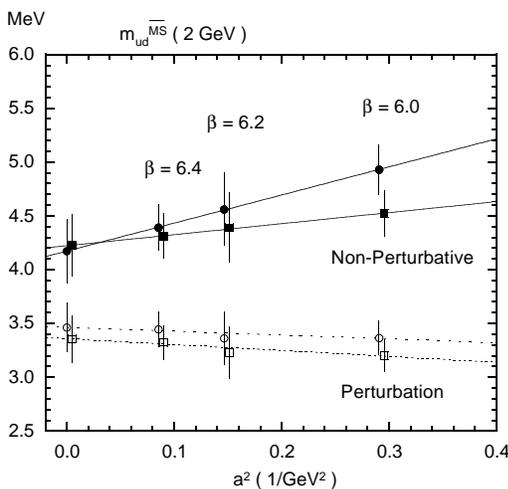


FIG. 4. The final results of the light quark mass $m_{ud}^{\overline{MS}}(2 \text{ GeV})$.

TABLE II. The nonperturbative (NP) and the perturbative (P) renormalization factors of the quark mass $Z_m(\mu, p, 1/a)$ at $\mu = 2$ GeV.

β	NP		P	
	$p = 1.8 \text{ GeV}$	$p = 2.6 \text{ GeV}$	$q^* = 1/a$	$q^* = \pi/a$
6.0	2.735(16)	2.509(18)	1.867	1.776
6.2	2.5451(91)	2.449(10)	1.877	1.800
6.4	2.385(12)	2.339(12)	1.871	1.804

For small momenta we consider that the momentum dependence arises from nonperturbative hadronization effects on the lattice and the higher order effects in continuum perturbation theory. It is very difficult to remove these effects from our results.

Toward large momenta, however, these effects are expected to disappear. The linear dependence on $(ap)^2$, which still remains, should arise from the discretization error on the lattice, i.e.,

$$Z_m(\mu, p, 1/a) = m^{\overline{MS}}(\mu)/m + (ap)^2 Z_H + O(a^4) \quad (8)$$

with the constant Z_H corresponding to the mixing to dimension 5 operators on the lattice.

This relation implies that, if we take a continuum extrapolation of $Z_m(\mu, p, 1/a)m$ at a fixed physical momentum p , the discretization error in Z_m is removed. This procedure also removes the a^2 discretization error in the lattice bare quark mass m itself, which reflects that in hadron masses.

The momentum p should be chosen in the region where the linear dependence on $(ap)^2$ is confirmed in our results. This region starts from a similar value of $p^2 \approx 3 \text{ GeV}^2$ for the three β values, and extends to $p^2 \approx 1.9/a^2$, the highest momentum measured. Hence we are able to use only a rather narrow range $3 < p^2 < 6.6 \text{ GeV}^2$, the upper bound dictated by the value of $1.9/a^2$ for the largest lattice spacing at $\beta = 6.0$.

In Fig. 4 we show the continuum extrapolation for the averaged up and down quark mass at $\mu = 2$ GeV. Filled circles are obtained for $p = 1.8 \text{ GeV}$ and squares for $p = 2.6 \text{ GeV}$ for which the value of Z_m is obtained by a linear fit in $(ap)^2$ employing the filled points in Fig. 3. The bare quark mass [10] is determined by a linear extrapolation of pseudoscalar meson mass squared in the Nambu-Goldstone channel ($\gamma_5 \otimes \xi_5$) and that of vector

TABLE III. Final results for $m_{ud}^{\overline{MS}}(2 \text{ GeV})$ obtained with the nonperturbative (NP) and the perturbative (P) renormalization factors. The lattice bare quark mass am_{ud} is also listed.

β	am_{ud} (10^{-4})	NP(MeV)		P(MeV)	
		$p = 1.8$ GeV	$p = 2.6$ GeV	$q^* = 1/a$	$q^* = \pi/a$
6.0	9.72(40)	4.93(23)	4.52(21)	3.37(15)	3.20(15)
6.2	6.86(45)	4.56(34)	4.39(32)	3.36(25)	3.23(24)
6.4	5.38(23)	4.39(21)	4.31(21)	3.45(16)	3.32(16)
$a^2 \rightarrow 0$		4.17(30)	4.23(29)	3.46(23)	3.36(22)

TABLE IV. The final results for the strange quark mass $m_s^{\overline{\text{MS}}}(2 \text{ GeV})$ obtained with K meson or ϕ meson mass to fix the bare strange quark masses am_s .

β	am_s (10^{-2})	NP(MeV)		P(MeV)	
		$p = 1.8$ GeV	$p = 2.6$ GeV	$q^* = 1/a$	$q^* = \pi/a$
K input					
6.0	2.44(10)	123.8(5.7)	113.6(5.3)	84.6(3.9)	80.4(3.7)
6.2	1.72(11)	114.4(3.3)	110.1(8.0)	84.5(6.2)	81.0(5.9)
6.4	1.350(57)	110.3(5.2)	108.2(5.1)	86.5(41)	83.4(3.9)
$a^2 \rightarrow 0$		104.7(7.4)	106.0(7.1)	86.9(5.6)	84.2(5.4)
ϕ input					
6.0	2.73(15)	138.5(8.2)	127.1(7.5)	94.6(5.4)	90.0(5.2)
6.2	2.04(23)	136(16)	131(15)	100(12)	96(11)
6.4	1.597(97)	130.4(8.4)	128(8.3)	102.3(6.6)	98.7(6.3)
$a^2 \rightarrow 0$		128(11)	129(12)	105.5(9.1)	102.2(8.7)

meson mass in the VT channel ($\overline{\gamma_i \otimes \xi_i}$) to the physical point of π and ρ meson masses.

We observe that the continuum extrapolation removes the momentum dependence of the quark mass at finite lattice spacings. Furthermore the values are substantially larger than those obtained with one-loop perturbation theory (open circles for $q^* = 1/a$ and squares for $q^* = \pi/a$). Making a linear extrapolation in a^2 , our final result in the continuum limit is

$$m_{ud}^{\overline{\text{MS}}}(2 \text{ GeV}) = 4.23(29) \text{ MeV}, \quad (9)$$

where we adopt the value for $p = 2.6 \text{ GeV}$ since this is the largest momentum accessible and the momentum dependence is negligible. The quoted error includes those arising from Z_m , the bare quark mass, and the lattice spacing. The value (9) is about 20% larger than the perturbative estimates: $3.46(23) \text{ MeV}$ for $q^* = 1/a$ and $3.36(22) \text{ MeV}$ for $q^* = \pi/a$. We collect the values of renormalization factor and quark masses in Tables II and III.

Applying our renormalization factor to the strange quark mass, we obtain

$$m_s^{\overline{\text{MS}}}(2 \text{ GeV}) = 106.0(7.1) \text{ MeV} \quad \text{for } m_K \quad (10)$$

$$= 129(12) \text{ MeV} \quad \text{for } m_\phi, \quad (11)$$

where we use K or ϕ meson mass to determine the bare strange mass. The results from perturbative estimation are given in Table IV.

The CP-PACS Collaboration recently reported the results [13] $m_{ud}^{\overline{\text{MS}}}(2 \text{ GeV}) = 4.6(2) \text{ MeV}$, $m_s^{\overline{\text{MS}}}(2 \text{ GeV}) = 115(2) \text{ MeV}(m_K)$, and $143(6) \text{ MeV}(m_\phi)$ from a large-scale precision simulation of hadron masses with the Wilson action. Our values are 10% smaller, which may be due to the use of a one-loop perturbative renormalization factor in the CP-PACS analysis.

This work is supported by the Supercomputer Project No. 32 (FY1998) of High Energy Accelerator Research Organization (KEK), and also in part by the Grants-in-Aid of the Ministry of Education (No. 08640404, No. 09304029, No. 10640246, No. 10640248, No. 10740107, and No. 10740125). S. K. and S. T. are supported by JSPS.

-
- [1] R. Gupta and T. Bhattacharya, Nucl. Phys. (Proc. Suppl.) **B63**, 95 (1998); R. D. Kenway, hep-lat/9810054.
 - [2] G. Martinelli *et al.*, Nucl. Phys. **B445**, 81 (1995).
 - [3] K. Jansen *et al.*, Nucl. Phys. **B372**, 275 (1996); S. Capitani, hep-lat/9810063.
 - [4] V. Giménez *et al.*, Nucl. Phys. **B540**, 473 (1999).
 - [5] D. Becirevic *et al.*, Phys. Lett. B **444**, 401 (1998).
 - [6] M. Golterman and J. Smit, Phys. Lett. **140B**, 392 (1984); A. Patel and S. Sharpe, Nucl. Phys. **B395**, 701 (1993); N. Ishizuka and Y. Shizawa, Phys. Rev. D **49**, 3519 (1994).
 - [7] D. Daniel and S. N. Sheard, Nucl. Phys. **B302**, 471 (1988).
 - [8] JLQCD Collaboration, S. Aoki *et al.*, hep-lat/9809124.
 - [9] For an earlier discussion, see G. Kilcup, R. Gupta, and S. Sharpe, Phys. Rev. D **57**, 1654 (1998).
 - [10] JLQCD Collaboration, S. Aoki *et al.*, Nucl. Phys. (Proc. Suppl.) **B53**, 209 (1997).
 - [11] G. P. Lepage and P. B. Mackenzie, Phys. Rev. D **48**, 2250 (1993).
 - [12] E. Franco and V. Lubicz, Nucl. Phys. **B531**, 641 (1998); T. van Ritbergen *et al.*, Phys. Lett. B **400**, 379 (1997); J. A. M. Vermaseren *et al.*, Phys. Lett. B **405**, 327 (1997); K. G. Chetyrkin, Phys. Lett. B **404**, 161 (1997).
 - [13] CP-PACS Collaboration, S. Aoki *et al.*, hep-lat/9809146.