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Variational Determination of the Stringtension and the Glueball Mass in (2 + 1) Dimensional U(1) Lattice Gauge Theory for all Values of the Coupling Constant

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Abstract. Using variational ansätze of the product type we calculate the massgap and the stringtension of $(2+1)$ dimensional $U(1)$ lattice gauge theory in the Hamiltonian formalism for all values of the coupling constant. In the strong coupling limit our results agree with high order strong coupling series. In the weak coupling limit both, glueball mass and stringtension vanish exponentially with the coupling constant.

1

Pure $U(1)$ lattice gauge theory in 3 space-time dimensions exhibits linear confinement [1–3] and has been shown to possess a nontrivial continuum limit [3]. In view of later applications to non-abelian gauge theories in higher dimensions this model is therefore well suited to serve as a laboratory for the development of simple approximation methods. Besides a correct description of the models properties, which are of nonperturbative nature in the weak coupling limit, these approximation methods should provide a transparent picture of the confinement mechanism. Additionally the quality of standard approximation methods like strong coupling expansions can be tested which might be instructive for their judgement even in more realistic models of confinement.

In this paper we use simple variational ansätze of the product type to calculate explicitly stringtension and glueballmass in the Hamiltonian formalism for all values of the coupling constant. Our results agree nicely with high order strong coupling series [4–5] and at least qualitatively with the exact result of Göpfert and Mack [3] and other approximate calculations [6, 7] in the weak coupling limit.

2

The Hamiltonian of $(2+1)$ dimensional $U(1)$ lattice gauge theory in the temporal gauge is given by

$$H=\frac{1}{2a}\sum_{\ell\in\Lambda} e^2E^2(l)+\frac{1}{2ae^2}\sum_{p\in\Lambda}(2-U(p)-U^+(p)) \tag{1}$$

with

$$U(p)=\prod_{\ell\in\partial p}U(\ell),\quad [E(\ell),U(\ell')]=\delta_{\ell'\ell}U(\ell');e^2=ag^2$$

denotes the dimensionless coupling constant, a the lattice constant and the sums extent over all links ℓ and unoriented plaquettes p of a finite, two-dimensional lattice Λ of $(2L)^2$ lattice sites. In the gauge invariant sector of the Hilbertspace $\mathscr{H}$ without charges the electric field operator $E(\ell)$ can be written as $E(\ell)=(\check{\nabla}F(p))(\ell)$ where $\check{\nabla}$ is the dual co-boundary operator. The operator $F(p)$, defined on plaquettes p is conjugate to the plaquette operator $U(p)$ in the sense of the commutation relation $[F(p),U(p')]=\delta_{p,p'}U(p')$. Using this new variable the Hamiltonian (1) becomes

$$H=\frac{1}{2a}\sum_{p\in\Lambda}\{-e^2F(p)(\nabla_+\cdot\nabla_-F)(p)+\frac{1}{e^2}(2-U(p)-U^+(p))\} \tag{2}$$

where $\nabla_+\nabla_-$ denotes the lattice Laplacian in terms of the discrete forward and backward lattice derivative ∇^j_+ and $\nabla^j_-, j=1,2$.

In the following we determine the properties of this Hamiltonian by variational methods. For this we consider as a first ansatz for the ground state

$$|\Omega_I\rangle=N_I\prod_{p\in\Lambda}\left\{\sum_{k=-\infty}^{\infty}e^{-(k^2/2\alpha_1)}U^k(p)\right\}|0\rangle \tag{3}$$

where the strong coupling vacuum $|0\rangle$ is defined by

$E(\ell)|0\rangle = F(p)|0\rangle = 0,$

and $\alpha_1(e^2)$ is considered as a variational parameter. This choice for the variational groundstate is motivated by the fact that

$$|\phi_0\rangle = N \cdot \sum_{k=-\infty}^{\infty} e^{-ag^2k^2} U^k(p)|0\rangle \tag{4}$$

is in the weak coupling limit, $e^2 = ag^2 \to 0$, an approximate solution for the groundstate of the single plaquette Hamiltonian [8]

$$a \cdot H_{sp} = 2e^2 F^2(p) + \frac{1}{2e^2}(2 - U(p) - U^+(p)). \tag{5}$$

In comparison to (3) we consider a second variational groundstate

$$|\Omega_{II}\rangle = N_{II} \exp\left\{\frac{\alpha_2}{4} \sum_{p \in A} (U(p) + U^+(p))\right\}|0\rangle \tag{6}$$

which has been used first for this model by Patkos [9] in the transfermatrix formalism.

3

Now we determine the two approximate groundstate energies E_0^I and E_0^{II}. A straightforward calculation gives for the expectation value of the Hamiltonian (2) with respect to the two states $|\Omega_I\rangle$ and $|\Omega_{II}\rangle$

$$E_0^I(\alpha_1) = \frac{1}{a}(2L)^2\left\{e^2\alpha_1\left(1 - \frac{A_1}{A_0}\right) + \frac{1}{e^2}\left(1 - \frac{B_0}{A_0} e^{(-1/4\alpha_1)}\right)\right\} \tag{7}$$

where

$$A_n = \left(-2\alpha_1 \frac{d}{d\alpha_1}\right)^n A_0, \quad A_0 = \sum_{m=-\infty}^{\infty} y^{m^2}$$

$$B_n = \left(-2\alpha_1 \frac{d}{d\alpha_1}\right)^n B_0, \quad B_0 = \sum_{m=-\infty}^{\infty} (-1)^m y^{m^2}$$

$y = \exp(-\pi\alpha_1)$

and

$$E_0^{II}(\alpha_2) = \frac{1}{2a}(2L)^2\left\{\left(\alpha_2 - \frac{2}{e^4}\right) I(\alpha_2) + \frac{2}{e^4}\right\} \tag{8}$$

respectively, where $I(x) = I_1(x)/I_0(x)$ denotes the ratio of the two modified Besselfunctions $I_0(x)$ and $I_1(x)$.

Minimization of $E_0^I(\alpha_1)$ and $E_0^{II}(\alpha_2)$ determines $\alpha_1(e^2)$ and $\alpha_2(e^2)$. Their functional dependence on the coupling constant is shown in Figs. 1 and 2. For their asymptotic behavior we find explicitly

$$\alpha_1(e^2)^{-1} \sim \begin{cases} 2\log 4e^4 & \text{for } e^2 \to \infty \\ e^2\sqrt{4\exp\frac{1}{4\alpha} + \mu^2} & \text{for } e^2 \to 0 \end{cases} \tag{9}$$

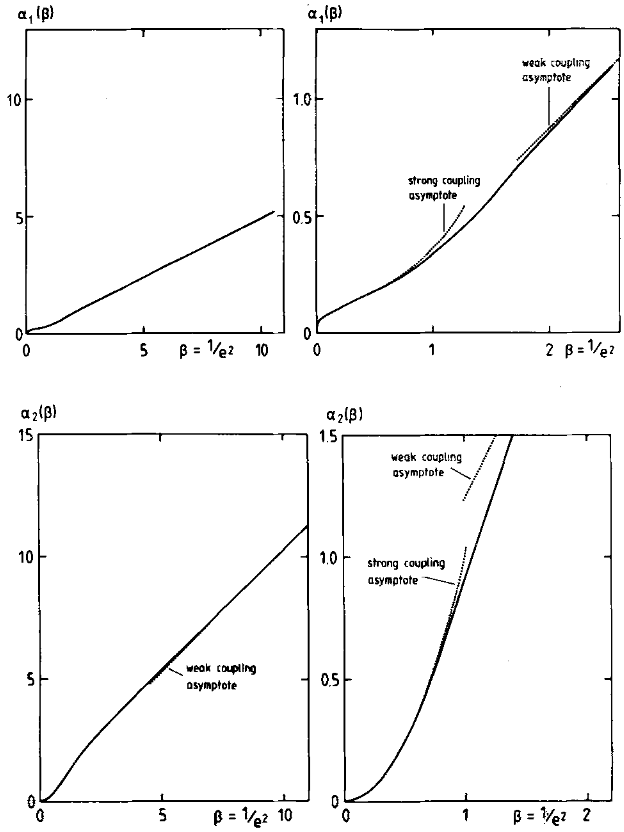

Fig. 1, 2. The variational parameters α_1, α_2 as function of the coupling constant $\beta = 1/e^2$ and their strong and weak coupling asymptotes

where

$$\mu^2 = \frac{4\pi^2}{e^4}(\pi^2 - 4)e^{-\pi^2\alpha}(1 + O(e^2) + O(y))$$

$$\alpha = \frac{1}{2e^2}\exp\left(-\frac{1}{8\alpha}\right) = \frac{1}{2e^2} - \frac{1}{8} - \frac{e^2}{64} + O(e^4)$$

and

$$\alpha_2(e^2) \sim \begin{cases} \frac{1}{e^4} & \text{for } e^2 \to \infty \\ \frac{1}{e^2} + \frac{63}{256} + O(e^2) & \text{for } e^2 \to 0 \end{cases} \tag{10}$$

Already from the asymptotic behavior of $\alpha_1(e^2)$ for small e^2 it can be seen that only the first ansatz (3) takes into account the non-perturbative tunnelling contributions which arise in the eigenproblem of the single plaquette Hamiltonian H_{sp} (5) and which are necessary for a correct description of the gauge theory in the weak coupling limit.

The results for the approximate energy densities of the groundstate $\lambda_0^{I,II} = E_0^{I,II}/(2L)^2$, are shown to-

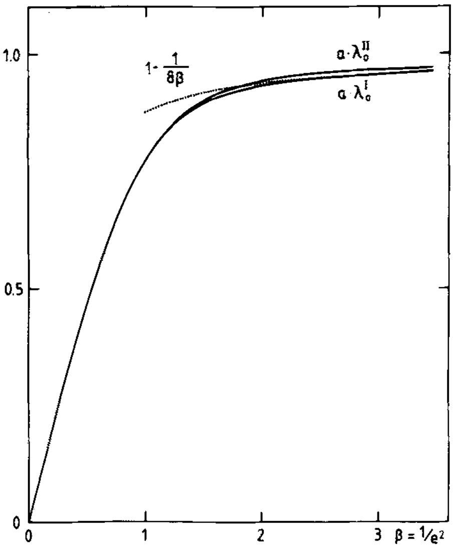

Fig. 3. The groundstate energy density $a\cdot\lambda_0^I, a\cdot\lambda_0^{II}$ of both trial states $|\Omega_I\rangle, |\Omega_{II}\rangle$ and their weak coupling asymptote $1-\frac{1}{8\beta}$

gether with their weak coupling expansion

$$a\cdot\lambda_0^{I,II} = 1 + \frac{1}{8}e^2 + O(e^4) \tag{11}$$

in Fig. 3. The first ansatz yields the lowest upper bound for the exact groundstate energy in the medium and weak coupling regime, while for strong couplings λ_0^I and λ_0^{II} are almost identical.

4

We determine now the massgap using both types of trial states. In the strong coupling region trial states of the product type yield the massgap directly. In the weak coupling regime, however, the results of an analogous variational calculation for a massive free scalar field have to be employed in order to separate the non-perturbative glueball mass.

According to (3) we choose for the trial state of first excited state

$$|\psi_1(\boldsymbol{\beta})\rangle_I = \frac{1}{2L}\sum_{\mathbf{p}\in\Lambda^*} e^{i\mathbf{p}\boldsymbol{\beta}}\psi_1(\mathbf{p})\prod_{\mathbf{q}\in\Lambda^*}{}'\psi_0(\mathbf{q})|0\rangle \tag{12}$$

where

$$\psi_0(\mathbf{q}) = \sum_{k=-\infty}^{\infty} e^{(-k^2/2\alpha_1)}U^k(\mathbf{q})|0\rangle$$

$$\psi_1(\mathbf{q}) = \sum_{k=-\infty}^{\infty} k\cdot e^{(-k^2/2\alpha_1)}U^k(\mathbf{q})|0\rangle \tag{13}$$

$$\beta_j = \frac{\pi}{L}n_j, n_j = -L,\ldots,L-1; j=1,2,$$

and $\mathbf{p}$ denotes a site of the dual lattice Λ^*. For the variational parameter $\alpha_1(e^2)$ we use the one determined before. This is justified for large values of L which we assume tacitly. For the energy $\varepsilon_1^I(\boldsymbol{\beta})$ of the first excited state we find explicitly

$$\varepsilon_1^I(\boldsymbol{\beta}) = \lambda_1^I - \lambda_0^I + (\mathbf{K}^2(\boldsymbol{\beta})-4)\cdot|M_I|^2 \tag{14}$$

where

$$\lambda_1^I = \frac{\langle\psi_1|H_{sp}|\psi_1\rangle}{\langle\psi_1|\psi_1\rangle}$$

$$= \frac{e^2\alpha_1}{a}\cdot\frac{3A_0-4A_1+A_2}{A_0-A_1} + \frac{1}{e^2a}$$

$$\cdot\left(1-\frac{\left(1-\frac{1}{2\alpha_1}\right)B_0-B_1}{A_0-A_1}e^{-(1/4\alpha_1)}\right)$$

$$|M_I|^2 = \frac{e^2}{a}|\langle\psi_1|F(p)|\psi_0\rangle|^2$$

$$= \frac{e\alpha_1}{2a}\left(1-\frac{A_1}{A_0}\right)$$

$$\mathbf{K}^2(\boldsymbol{\beta}) = \sum_{j=1,2} 4\sin^2\frac{1}{2}\beta_j; 0 \leqq \mathbf{K}^2(\boldsymbol{\beta}) \leqq 8.$$

The variational massgap $\varepsilon_1^I(\mathbf{0})$ is shown in Fig. 4 in comparison with its 8th order strong coupling series $\varepsilon_1^{sc}(\mathbf{0})$ [4]. The agreement is better than 7% for values of the coupling constant $1/e^2 \leqq 0.9$ for which the strong coupling series seems to converge.

As mentioned before a detailed analysis of $\varepsilon_1^I(\mathbf{0})$ is necessary to reveal the glueball mass in the weak coupling limit where the gauge theory is described by a free massive scalar field theory [3] whose mass vanishes for $e^2 \to 0$ and constant lattice spacing a. As a first check of our result (14) we observe that both,

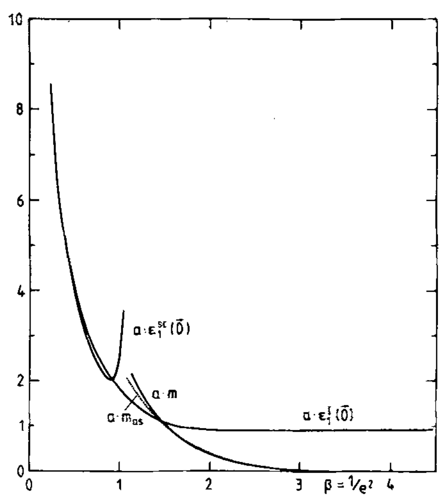

Fig. 4. The massgap calculated by a) strong coupling series [4], $a\cdot\varepsilon_1^{sc}(\mathbf{0})$, b) variational ansatz $|\psi_1\rangle_I, a\cdot\varepsilon_1^I(\mathbf{0})$, c) variational ansatz $|\psi_1\rangle_I$ and comparison with the free massive scalar field, $a\cdot m$, and its asymptote $a\cdot m_{as}$ as a function of $\beta = 1/e^2$

the variational massgap of the pure gauge theory $\varepsilon_1^I(\mathbf{0})$ and the one of the scalarfield, $\varepsilon_1^{\rm sf}(\mathbf{0})$, of mass $m_{\rm sf}=\mu_{\rm sf}/a$

$$\varepsilon_1^{\rm sf}(\boldsymbol{\beta})=\sqrt{m_{\rm sf}^2+\frac{4}{a^2}}\cdot\left(1+\frac{1}{2}\,\frac{\frac{1}{a^2}\mathbf{K}^2(\boldsymbol{\beta})-\frac{4}{a^2}}{m_{\rm sf}^2+\frac{4}{a^2}}\right) \tag{15}$$

approach $1/a$ in the limit $e^2\to 0$, g fixed, and $m_{\rm sf}^2\to 0$, respectively. Equation (15) has been obtained with help of variational ansätze analogous to (3),(13). It is identical to the result of a first order perturbation calculation where the Hamiltonian $H_{\rm sf}$ of the scalarfield $\phi(\mathbf{p})$ is split into two parts

$$H_{\rm sf}=H_{\rm sf}^0+H_{\rm sf}^I$$

$$H_{\rm sf}^0=\frac{1}{2a}\sum_{\mathbf{p}\in\Lambda^*}\{\pi^2(\mathbf{p})+(4+\mu_{\rm sf}^2)\phi^2(\mathbf{p})\}$$

$$H_{\rm sf}^I=\frac{1}{2a}\sum_{\mathbf{p}\in\Lambda^*}\{-\phi(\mathbf{p})((\nabla_+\cdot\nabla_-+4)\phi)(\mathbf{p})\} \tag{16}$$

and $H_{\rm sf}^I$ is treated as perturbation. It has been verified by Schiff [10] up to third order that higher order perturbation calculations improve the result (15) for $\varepsilon_1^{\rm sf}(\boldsymbol{\beta})$ and yield the first terms of the convergent power expansion

$$\varepsilon_1(\boldsymbol{\beta})=\sqrt{m_{\rm sf}^2+\frac{1}{a^2}\mathbf{K}^2(\boldsymbol{\beta})}$$

$$=\sqrt{m_{\rm sf}^2+\frac{4}{a^2}}\cdot(1+\tfrac{1}{2}\Delta(\boldsymbol{\beta})-\tfrac{1}{8}\Delta^2(\boldsymbol{\beta})+\cdots)$$

$$\Delta(\boldsymbol{\beta})=\left(\frac{1}{a^2}\mathbf{K}^2(\boldsymbol{\beta})-\frac{4}{a^2}\right)\cdot\left(m_{\rm sf}^2+\frac{4}{a^2}\right)^{-1} \tag{17}$$

of the relativistic energy of a massive scalarfield of lattice momentum $1/a\mathbf{K}(\boldsymbol{\beta}),(\mathbf{K}_j(\boldsymbol{\beta}_j)=2\sin 1/2\boldsymbol{\beta}_j$; $j=1,2)$. The expansion in (17) is performed in terms of the relative deviation $\Delta(\boldsymbol{\beta})$ of the square of the lattice momentum from half of its maximal value $1/2[1/a^2\mathbf{K}^2(\boldsymbol{\beta})]_{\max}=4/a^2$.

The comparison of (14,15) reveals the nonperturbative glueball mass in the weak coupling limit. We find

$$m^2(e^2)=\frac{\lambda_1^I-\lambda_0^I}{2|M_I|^2}-\frac{4}{a^2} \tag{18}$$

$$\underset{e^2\to 0}{\sim}\frac{4\pi^2}{a^2e^4}(\pi^2-4)e^{-\pi^2\alpha}(1+O(e^2)+O(y))$$

The same asymptotic behavior for $m^2(e^2)$ is obtained by the comparison of the variational groundstate for the massive scalarfield (16) and the variational groundstate of the gauge theory (3), (9) and also from the heavy $q\bar{q}$-state in the next section. Figure 4 shows the glueball mass and its weak coupling asymptote in comparison with $\varepsilon_1^I(\mathbf{0})$ and $\varepsilon_1^{\rm sc}(\mathbf{0})$, $m(e^2)$ matches its weak coupling asymptote at relatively large couplings

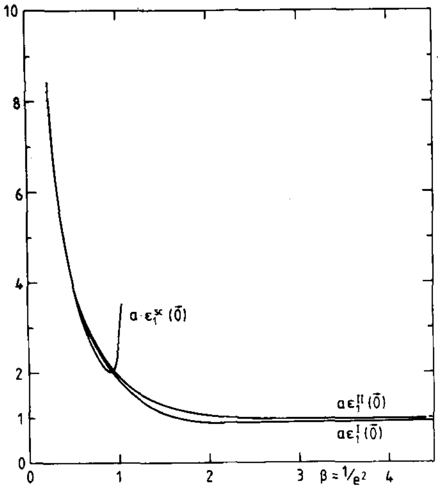

Fig. 5. The massgaps $a\cdot\varepsilon_1^I(\mathbf{0})$ and $a\cdot\varepsilon_1^{II}(\mathbf{0})$ in comparison with the strong coupling expansion $a\cdot\varepsilon_1^{\rm sc}(\mathbf{0})$

$1/e^2\gtrsim 1.5$ which indicates an early continuum behavior.

A similar result for the asymptotic behavior of the glueball mass in the weak coupling region has been obtained independently by Suranyi [7] where the different exponential dependence on the coupling constant is due to his variational ansatz which is not of the product type.

The result $\varepsilon_1^{II}(\mathbf{0})$ of a variational calculation of the massgap using an ansatz of the second type, (12) but

$$\psi_0(\mathbf{q})=\exp\left\{\frac{\alpha_2}{4}(U(\mathbf{q})+U^+(\mathbf{q}))\right\}$$

$$\psi_1(\mathbf{q})=(U(\mathbf{q})-U^+(\mathbf{q}))\exp\left\{\frac{\alpha_2}{4}(U(\mathbf{q})+U^+(\mathbf{q}))\right\} \tag{19}$$

is shown in Fig. 5 in comparison with $\varepsilon_1^I(\mathbf{0})$ and $\varepsilon_1^{\rm sc}(\mathbf{0})$. They all agree nicely for strong couplings. In the weak coupling regime, however, the second ansatz does not allow the determination of a nonperturbative, exponentially small glueball mass. This failure of the second ansatz clearly shows the necessity of the proper treatment of arbitrary high plaquette excitations. These excitations dominate the first excited states of (13), (19) in the weak coupling limit and are indispensable for the correct description of the tunnelling contributions and thus of the glueball mass and the stringtension for small couplings.

The importance of plaquette excitations of high order also has been realized recently by Hamer et al. [11] in a finite lattice Hamiltonian calculation.

5

For the calculation of the heavy quark potential we add to the trial groundstate two static sources, $q^+(\mathbf{r})$ and $q(\mathbf{s})$, which are connected by the parallel transport along a path $\mathscr{C}(\mathbf{r},\mathbf{s})$. We use the two trial states

$$|S_I[\mathscr{C}(\mathbf{r},\mathbf{s})]\rangle = \tilde{N}_I \hat{\mathscr{C}}(\mathbf{r},\mathbf{s}) \prod_{p\in A} \cdot \left\{ \sum_{k=-\infty}^{\infty} e^{-(k-\delta(p))^2/2\alpha_1} U^k(p) \right\} |0\rangle \qquad (20)$$

and

$$|S_{II}[\mathscr{C}(\mathbf{r},\mathbf{s})]\rangle = N_{II} \hat{\mathscr{C}}(\mathbf{r},\mathbf{s}) \cdot \exp\left\{ \frac{\alpha_2}{4} \sum_{p\in A} (\alpha_2 + \varepsilon(p)) U(p) + (\alpha_2 - \varepsilon(p)) U^+(p)) \right\} |0\rangle$$

$$\text{where } \hat{\mathscr{C}}(\mathbf{r},\mathbf{s}) = q^+(\mathbf{r}) \prod_{\ell\in\mathscr{C}(\mathbf{r},\mathbf{s})} U(\ell) q(\mathbf{s}). \qquad (21)$$

The new, real and space-dependent variational parameters $\varepsilon(p)$ and $\delta(p)$ allow for a broadening of the string $\hat{\mathscr{C}}(r,s)$ into a tube of electric flux since now the expectation value of the electric field operator in general does not vanish on links away from the string, rather

$$\langle S_I | E(\ell) | S_I \rangle = (\check{\nabla}\eta)(\ell) \quad \ell \notin \mathscr{C}(\mathbf{r},\mathbf{s})$$

where

$$\eta(p) = \delta(p) + O\left(\frac{1}{e^2} \exp - \frac{\pi^2}{2e^2} \right) \qquad (22)$$

and

$$\langle S_{II} | E(\ell) | S_{II} \rangle = (\check{\nabla}\phi)(\ell) \quad \ell \notin \mathscr{C}(\mathbf{r},\mathbf{s}) \qquad (23)$$

where

$$\phi(p) = \tfrac{1}{2}\varepsilon(p) \cdot I(\alpha_2)$$

respectively.

The variational energy $E_s^{II}[\mathscr{C}(\mathbf{r},\mathbf{s})]$ of the second trial state (21) is given by the energy of a classical, euclidean free scalar field of mass $m_2(e^2)$ which is coupled to the string $\hat{\mathscr{C}}(\mathbf{r},\mathbf{s})$ as an external source. The string tension for a straight string can be calculated for all values of the coupling constant with the result

$$T_2(e^2) = \frac{e^2}{2a} \frac{m_2}{\sqrt{m_2^2 + \frac{4}{a^2}}} \qquad (24)$$

$$\text{where} \quad m_2^2 = \frac{-4e^2}{a^2(2 - e^4\alpha_2) I(\alpha_2)}.$$

It stays positive for all values of $e^2 > 0$, matches its strong coupling series [5] for large e^2 and vanishes perturbatively in the weak coupling region. $T_2(e^2)$ is shown together with the 10th order strong coupling series $T_{sc}(e^2)$ of [5] in Fig. 6. In the weak coupling region $T_2(e^2)$ is not exponentially small as anticipated [1–3]. This behavior, however, is improved by the trial state $|S_I\rangle$ of (20). Its variational energy $E_s^I(\mathbf{r},\mathbf{s})$ can be minimized in the weak coupling region for infinitely separated sources located on the same main axis of the lattice, where it essentially becomes the classical energy of the kink of the inhomogeneous, one-

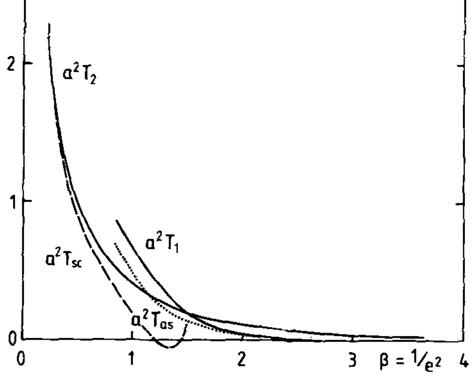

Fig. 6. The stringtension $a^2 T_1$ and $a^2 T_2$ and its strong coupling expansion $a^2 T_{sc}$ of 10th order of [5]

dimensional, euclidean Sine–Gordon equation on the lattice. This leads to the stringtension

$$T_1(e^2) = \frac{2e^2}{a} \frac{m}{\pi} + O(y) \qquad (25)$$

where m is given by (18). This is in agreement with similar results of Suranyi [7] and Heller [8]. The flux tube connecting both charges is of finite width in both trial states (20), (21) since both expectation values of the electric field operator (22), (23) decrease exponentially in the transverse direction to the $q\bar{q}$-axis. The width σ^2 determined by the measure of Lüscher et al. [12] for the first trial state $|S_I\rangle$ in the weak coupling limit is the same as the one of [7]

$$\sigma^2 = \frac{\pi^2}{12} \cdot \frac{1}{m^2}$$

though a different trial state was used here. The glueball mass which follows from the exponential fall-off of $\langle S_I | E(\ell) | S_I \rangle$ agrees with the one obtained before (18)

Both types of trial states yield in the weak coupling limit similar results for the ratio of Göpfert and Mack [3]

$$R = \frac{T}{m_G} \cdot \frac{4\pi^2}{e^2} \cdot a$$

We find $R_I \underset{e^2\to 0}{\to} 8$, which is the result of Göpfert and Mack and $R_{II} \underset{e^2\to 0}{\to} \pi^2$.

6. Conclusion

As a result of our study of (2 + 1) dimensional $U(1)$ lattice gauge theory we conclude that simple variational ansätze of the product type are sufficient to describe the known properties of the model for all values of the coupling constant. While in the strong coupling limit both types of trial states used here led to similar results, the correct treatment of high plaquette excitations proved to be essential in the medium and weak coupling region for an at least qualitatively correct description of the nonperturbative glueballmass and stringtension. The extension of the performed analysis to non-abelian gauge theories should improve the existing variational results [13, 14] in the weak coupling regime. The extension to higher space-time dimensions would be of great interest.

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