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Dynamical W mass?

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As long as a Higgs boson is not observed, the design of alternatives for electroweak symmetry breaking remains of interest. The question addressed here is whether there are possibly dynamical mechanisms, which deconfine SU(2) at zero temperature and generate a massive vector boson triplet. Results for a model with joint local U(2) transformations of SU(2) and U(1) vector fields are presented in a limit, which does not involve any unobserved fields.

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1. Introduction

In Euclidean field theory notation the action of the electroweak gauge part of the standard model reads

$$S = \int d^4x L^{\rm ew}\,, \quad L^{\rm ew} = -\frac{1}{4}F^{\rm em}_{\mu\nu}F^{\rm em}_{\mu\nu} - \frac{1}{2}{\rm Tr}F^b_{\mu\nu}F^b_{\mu\nu}\,, \tag{1.1}$$

$$F^{\rm em}_{\mu\nu} = \partial_\mu a_\nu - \partial_\nu a_\mu\,, \quad F^b_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu + ig_b\left[B_\mu, B_\nu\right]\,, \tag{1.2}$$

where a'_μ are U(1) and B_μ are SU(2) gauge fields.

Typical textbook introductions of the standard model emphasize at this point that the theory contains four massless gauge bosons and introduce the Higgs mechanism as a vehicle to modify the theory so that only one gauge boson, the photon, stays massless. Such presentations reflect that the introduction of the Higgs particle in electroweak interactions [1] preceded our non-perturbative understanding of non-Abelian gauge theories. In fact, massless gluons are not in the physical spectrum of (1.1).

The self interaction due to the commutator (1.2) generates dynamically a non-perturbative mass gap, and the SU(2) spectrum consists of massive glueballs. The lightest glueball can be used to set a mass scale. Choosing for it, e.g., 80 GeV and coupling fermions is admissible. This does not constitute an ansatz for an electroweak theory by two main reasons: The SU(2) glueball spectrum [2] is not what is wanted (e.g., masses of spin 0 and 2 states are lower than for spin 1), and, fermions would be confined, which is not the case. Coupling a Higgs field in the usual way causes a deconfining phase transition, so that fermions are liberated, a photon stays massless and glueballs break up into elementary massive vector bosons. Such a confinement-Higgs transition has indeed been observed in pioneering lattice gauge theory (LGT) investigations [3].

This mass generation for the W boson through spontaneous symmetry breaking is explicit in perturbation theory. Are there possibly alternative dynamical mechanisms, which do not involve new physics? The model considered here is motivated by the deconfining phase transition of pure U(1) LGT with the Wilson action [4], which is demonstrated in Fig. 1 with Polyakov loop scatter plots for coupling constants in the confined ($\beta_e = 0.9$) and Coulomb ($\beta_e = 1.1$) phase. In the Coulomb phase the effective potential of the Polyakov loop is similar to that of a Higgs field.

2. SU(2) alignment

Motivated by the behavior of the U(1) Polyakov loop in the Coulomb phase, we add to the U(1) and SU(2) Wilson actions

$$S^{\rm add} = \sum_{\mu\nu} S^{\rm add}_{\mu\nu}\,, \quad S^{\rm add}_{\mu\nu} = \frac{\lambda}{2}{\rm Re\,Tr}\left[U_\mu(x)V_\nu(x+\hat\mu a)U^*_\mu(x+\hat\nu a)V^*_\nu(x)\right] \tag{2.1}$$

with $U \in U(1)$ taken as diagonal 2×2 matrices and $V \in SU(2)$. For aligned U(1) matrices the SU(2) matrices become aligned too and for large enough λ one may expect a SU(2) deconfining phase transition due to breaking of the Z_2 center symmetry.

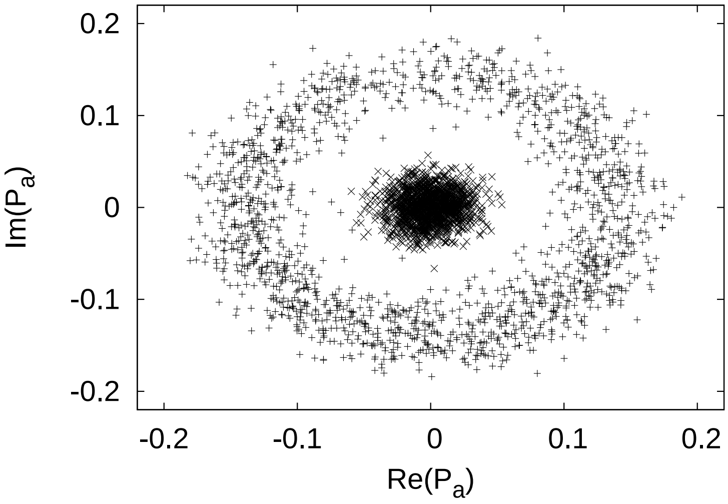

Figure 1: Scatter plot for U(1) Polyakov loops with the Wilson action on a 12^4 lattice at $\beta_e = 0.9$ (center) and $\beta_e = 1.1$ (ring), $\beta_e = 1/g_e^2$.

This interaction can be obtained in the London limit ($\kappa \to \infty$) from the following gauge invariant expression:

$$S^{add}_{\mu\nu} \;=\; \kappa\,\mathrm{Tr}\left[(\Phi^+\Phi - \mathbf{1})^2\right]\} + \frac{\lambda}{2}\mathrm{Re}\,\mathrm{Tr}\left\{U_\mu(x)\,U^*_\mu(x+\hat{\nu}a)\right. \qquad (2.2)$$
$$\left.[\Phi^+(x+\hat{\mu}a)V_\nu(x+\hat{\mu}a)\Phi(x+\hat{\mu}a+\hat{\nu}a)]\,[\Phi^+(x)V_\nu(x)\Phi(x+\hat{\nu}a)]^+\right\}$$

where Φ is a 2×2 matrix scalar field that is charged with respect to $U(1)$ and $SU(2)$. The gauge transformations are $\Phi \to e^{-i\alpha}g\,\Phi$, where $g \in SU(2)$, $e^{i\alpha} \in U(1)$. The vacuum value of Φ is a pure gauge $\Phi = e^{-i\alpha}g$. To obtain (2.1) we perform the the London limit for the potential and fix the gauge to $\Phi = \mathbf{1}$. Similar results are expected from simulations at sufficiently large finite κ values.

In the classical continuum limit $a \to 0$ the a^4 contributions of $S^{\rm add}_{\mu\nu}$ (after gauge fixing) give

$$L^{\rm add} \;=\; -\frac{\lambda}{2}\,\mathrm{Tr}\left(F^{\rm add}_{\mu\nu}F^{\rm add}_{\mu\nu}\right)\,, \quad F^{\rm add}_{\mu\nu} \;=\; g_a\partial_\mu A_\nu - g_b\partial_\nu B_\mu \qquad (2.3)$$

where $A_\mu = a_\mu\mathbf{1}$ is the photon and B_μ the gluon field. The interaction leads to similarities with the SU(2) Higgs model. However, there is no explicit mass term $\sim B_\mu B_\mu$ which would be obtained by applying the London limit to conventional Higgs coupling.

Our Monte Carlo procedure proposes the usual U(1) and SU(2) changes. For the update of a SU(2) matrix $V_\mu(n)$ we need the contribution to the action, which comes from the summed up staples (using now lattice units $a=1$)

$$V_{\sqcup,\mu}(n) \;=\; \frac{\beta_b}{2}\sum_{\nu\neq\mu}\left[V_\nu(n+\hat{\mu})\,V^*_\mu(n+\hat{\nu})\,V^*_\nu(n) + V^*_\nu(n+\hat{\mu}-\hat{\nu})\,V^*_\mu(n-\hat{\nu})\,V_\nu(n-\hat{\nu})\right]$$
$$+\;\frac{\lambda}{2}\sum_{\nu}\left[U_\nu(n+\hat{\mu})\,V^*_\mu(n+\hat{\nu})\,U^*_\nu(n) + U^*_\nu(n+\hat{\mu}-\hat{\nu})\,V^*_\mu(n-\hat{\nu})\,U_\nu(n-\hat{\nu})\right]\,,$$

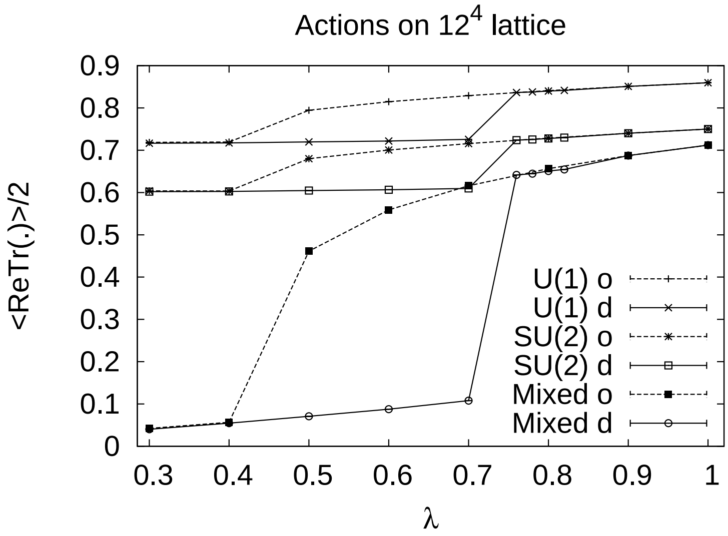

Figure 2: Plaquette expectation values on a 12^4 lattice as function of λ (ordered o and disordered d starts).

and correspondingly for the U(1) matrices $U_\mu(n)$. This is well suited for updates with a biased Metropolis-heatbath algorithm [5], which gives acceptance rates larger than 95% for U(1) as well as for SU(2) updates in the range of parameters considered.

3. Numerical results

Simulations reported here are at $\beta_e = 1.1$ in the Coulomb phase of pure U(1) LGT and $\beta_b = 2.3$ in the scaling region of pure SU(2) LGT. In λ a strong first-order transition is found as shown in Fig. 2 for plaquette expectation values, which are normalized to one for unit matrices.

SU(2) and U(1) string tensions from Creutz ratios [6] of Wilson loops up to size 5×5 as shown in Fig. 3, as well as Polyakov loop histograms (not shown), support a SU(2) deconfining phase transition, while U(1) is deconfined on both sides with a discontinuity in the Coulomb potential.

Relying on the dispersion relation, the photon mass is estimated [7] from fits to correlation functions for photon energy E_{k_1}, $k_1 = 2\pi/N$ estimates on N^3N_t, $N_t \gg N$ lattices, which are shown in Fig. 4. Their analysis yields

$$m^2_{\rm photon} = E^2_{k_1} - 4\,\sin^2(k_1/2) \to 0 \tag{3.1}$$

with increasing N on both sides of the transition.

Other mass spectrum estimates are obtained from fits to zero-momentum correlations functions. Glueball correlations are very noisy, best for 0^+, and signals are only followed up to distance 2 in the disordered phase and even worse (higher masses) in the ordered phase.

The mass m_W of the vector boson triplet is estimated from correlations of the operator

$$V_{i,\mu}(x) \;=\; -i\,{\rm Tr}\left[\tau_i V_\mu(x)\right] . \tag{3.2}$$

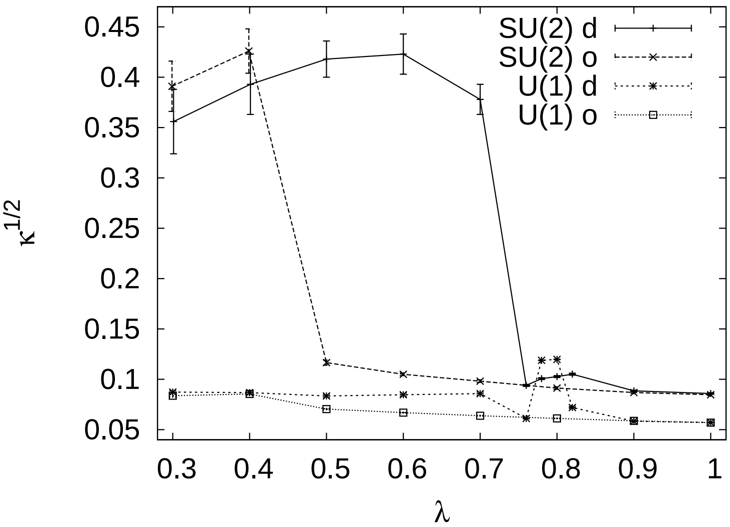

Figure 3: SU(2) and U(1) string tensions on a 12^4 lattice as function of λ (disordered d and ordered o starts).

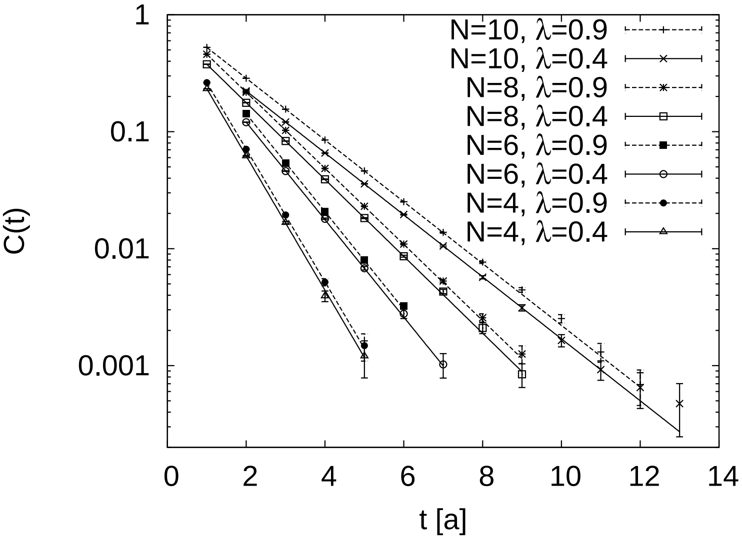

Figure 4: Data and fits for photon energy E_{k_1} correlation functions. The up-down order of the curves agrees with that of the labeling.

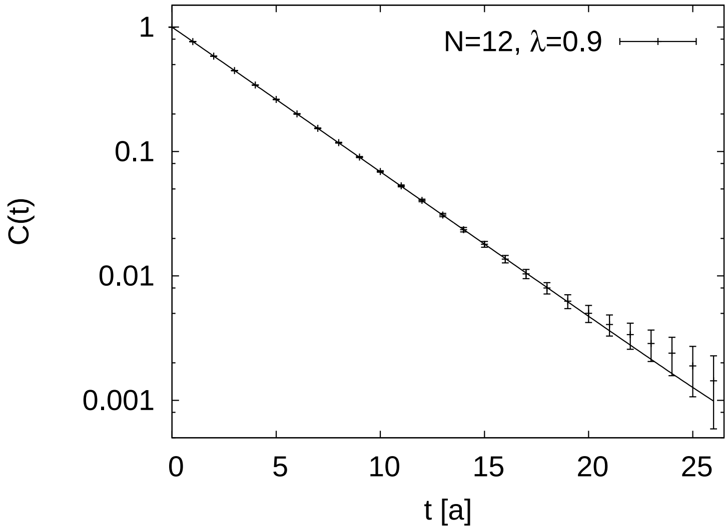

Figure 5: Momentum zero correlation function and fit for m_W mass estimate at $\lambda = 0.9$ on a $12^3 64$ lattice.

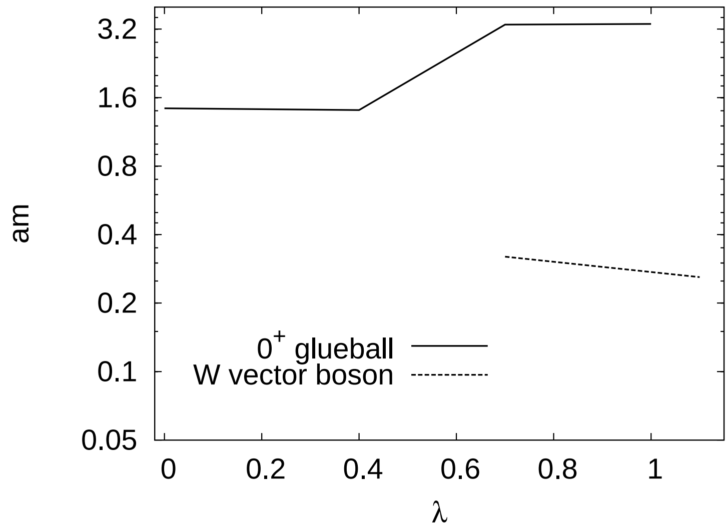

Figure 6: Sketch of the glueball and vector boson mass spectrum.

They are zero in the disordered phase and strong in the ordered phase. For our largest lattice Fig. 5 shows them at $\lambda = 0.9$ (ordered start). The signal is beautifully strong and can be followed beyond distance $t = 20\,a$. A sketch of the glueball and W mass spectrum as function of λ is given in Fig. 6.

4. Summary and conclusions

To the extent that similar results hold also for finite (sufficiently large) values of κ, we have constructed a gauge-invariant U(1)$\otimes$SU(2) theory, which exhibits a zero-temperature deconfining transition and generates a W boson mass dynamically. However, by dimensional reasons this theory with scalar bosons has presumably no physical quantum continuum limit.

In the London limit the U(1)⊗SU(2) scalar matrix field Φ becomes unphysical, because it takes its vacuum value and does not fluctuate, while its gauge transformations survive. In the present model they can be absorbed by extending the gauge transformations of the U(1) and SU(2) vector fields into a local U(2)=U(1)⊗SU(2) symmetry

$$U_\mu(x) \to e^{-i\alpha(x)} g(x) U_\mu(x) g^{-1}(x+\hat{\mu}a) e^{i\alpha(x+\hat{\mu}a)} ,$$
$$V_\mu(x) \to e^{-i\alpha(x)} g(x) V_\mu(x) g^{-1}(x+\hat{\mu}a) e^{i\alpha(x+\hat{\mu}a)} ,$$

which was first proposed in [8, 9].

In this way all remnants of the scalar field disappear without destroying invariance of the action, so that a local U(2) invariance of (fermionic) matter fields can be kept. After elimination of the scalar bosons the additional action term is now (as desired) dimensionless and it remains to clarify whether this allows for a quantum continuum limit. The situation is kind of opposite to the SU(2) Higgs model, where the interaction term is dimensionless as long as the scalar boson is involved, but acquires a dimension in the London limit.

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