Chapter 14

Hamiltonian Formulation of Local Symmetry

For several purposes, it is sometimes preferable to work on a lattice with different lattice spacings in different directions, as we have done in analyzing the plaquette. We can write the action as a sum over sites with the appropriate space-time volume,

$$\int d^4x \rightarrow a^x a^y a^z a^t \sum_n$$

$$U_{n,n+\mu} = e^{i a^\mu A_\mu}$$ on each link

$$F_{\mu\nu}^2 \rightarrow - \sum_{\mu\nu} \frac{1}{g^2 a^2 a^\nu} \text{Tr} \left( U_{\mu\nu} + U^\dagger_{\mu\nu} - 2 \cdot \mathbb{I} \right)$$ for each plaquette

and so

$$S = a^x a^y a^z a^t \sum_n \left\{ \frac{1}{4g^2} \sum_{\mu\nu} \frac{1}{a^\mu a^\nu} \text{Tr} \left( U_{\mu\nu} + U^\dagger_{\mu\nu} - 2 \cdot \mathbb{I} \right) (-1)^{\delta_{\mu\nu} + \delta_{\nu\nu}} + i \sum_{\mu} \frac{1}{a^\mu} \bar{\psi}_{n+\mu} \gamma^\mu U_{n,n+\mu} \psi_n - m \bar{\psi} \psi \right\}.$$

We now turn to a discussion of the Hamiltonian formulation, as well as the Lagrangian. This will give us a quantum mechanics in the usual language, but of a rather unusual form, because the kinetic energy term will not be the usual $\sum p_i^2/2m$, because the degrees of freedom, the group elements on each link, do not live in a Euclidean space, but instead live on the group manifold. The way the derivative operators enter the Hamiltonian will lead us to a metric and a measure on the group manifold.

We will derive the Hamiltonian formulation of the lattice gauge theory by considering the $a^t \rightarrow 0$ limit. First we must use gauge invariance to remove some of the degrees of freedom.

Quantum mechanics is an integral of the action over all possible configurations of the physical degrees of freedom. Gauge degrees of freedom have no effect on the action, and are therefore not physical degrees of freedom. Given any configuration $U, \psi$ on all the links and sites, it is possible to find a new configuration, gauge transformed, satisfying some imposed gauge requirements but representing the same physical configuration. In our case, we pick a particular time, say $t = 0$. We now perform gauge transformations at each site with $t \neq 0$ in order to make all $U_{n,n+i} = \mathbb{I}$. This can be done iteratively. For example, at $t = a^t$, at site $n = \vec{r}, a^t$ we examine $U_{(\vec{r},0),\hat{r}} = e^{i \omega \vec{r}}$, and then make a gauge transformation at the site $(\vec{r}, a^t)$ with $\omega \vec{r}$, with no change at $(\vec{r}, 0)$, so the new $U$ on the timelike link is

$$U'_{(\vec{r},0),\hat{r}} = e^{-i \omega \vec{r}} e^{i \omega \vec{r}} = \mathbb{I}.$$

This set of transformations changes all the spacelike $U$’s at $t = a^t$ to new values, but it does not affect the integral over all such $U$’s, provided the weight of integration is group invariant. We shall say more about this later.

Reiterating this step forwards in time to $t = \infty$, and also going backwards from $t = 0$, we wind up with a configuration of $\psi$’s at $\psi_{(\vec{r}, a^t)}$, with no change at $(\vec{r}, 0)$. The sum over times, times $a^t$, is just $\int dt$, so $S = \int dt L$ as usual, with

$$L = \sum_{n_x,n_y,n_z} a^3 \left\{ \frac{1}{2g^2 a^2 a^t} \sum_j \left( \text{Tr} U_{p_j} - \mathbb{I} + \text{h.c.} \right) + \frac{1}{4g^2 a^2} \sum_{ij} \text{Tr} \left( U_{p_i} + U_{p_i}^\dagger - 2 \cdot \mathbb{I} \right) \right\} + \text{fermion terms}.$$

The first term is the sum over plaquettes in a time-space plane, while the
second is in a space-space plane (at a fixed time). The first of these involves
\[ U_{p_{ij}} = U_{n^\dagger,0}^\dagger U_{n^\dagger,0} U_{n,0}. \]

Expand
\[ U_{(\bar{a},t+a^\prime),j} = U_{(\bar{a},t),j} + a^\prime \frac{\partial U_{(\bar{a},t),j}}{\partial t} + \frac{1}{2} a^\prime \frac{\partial^2 U}{\partial t^2}, \]
so
\[ U_{(\bar{a},t+a^\prime),j} U_{(\bar{a},t),j} - \mathbb{I} = a^\prime \left( \frac{\partial U^{-1}_{n^\dagger,j}}{\partial t} \right) U_{n^\dagger,j} + \frac{a^{\prime 2}}{2} \frac{\partial^2 U^{-1}_{n^\dagger,j}}{\partial t^2} U_{n^\dagger,j}. \]

Adding the hermitean conjugate gives
\[ a^\prime \left( \frac{\partial U^{-1}_{n^\dagger,j}}{\partial t} U_{n^\dagger,j} + U_{n^\dagger,j} \frac{\partial U^{-1}_{n^\dagger,j}}{\partial t} \right) + \frac{a^{\prime 2}}{2} \left( \frac{\partial^2 U^{-1}}{\partial t^2} U + U^{-1} \frac{\partial^2 U}{\partial t^2} \right). \]
The first term is \( a^\prime \frac{\partial}{\partial t} (U^{-1} U) = a^\prime \left( \frac{\partial}{\partial t} \right) \mathbb{I} = 0, \) while from the second we may subtract
\[ \frac{a^{\prime 2}}{2} \frac{\partial^2}{\partial t^2} 0 = \frac{a^{\prime 2}}{2} \frac{\partial}{\partial t} \left( \frac{\partial U^{-1}}{\partial t} U + U^{-1} \frac{\partial U}{\partial t} \right) = \frac{a^{\prime 2}}{2} \left( \frac{\partial^2 U^{-1}}{\partial t^2} + 2 \frac{\partial U^{-1}}{\partial t} \frac{\partial U}{\partial t} + U^{-1} \frac{\partial^2 U}{\partial t^2} \right), \]
which leaves
\[ \text{Tr} \left( U_{p_{ij}} - 1 \right) + \text{h.c.} = \text{Tr} \left( U_{t+a^\prime}^\dagger U_t - \mathbb{I} \right) + \text{h.c.} = -a^{\prime 2} \text{Tr} \left( \frac{\partial U^{-1}}{\partial t} \frac{\partial U}{\partial t} \right). \]
We can improve this by using our expression for zero again, \( \frac{\partial U^{-1}}{\partial t} U = -U^{-1} \frac{\partial U}{\partial t}, \) so \( \frac{\partial U^{-1}}{\partial t} = -U^{-1} \frac{\partial U}{\partial t} U^{-1} \) and
\[ \text{Tr} \left( U_{t+a^\prime}^\dagger U_t - \mathbb{I} \right) + \text{h.c.} = a^{\prime 2} \text{Tr} \left( U^{-1} \dot{U} U^{-1} \right), \]
and
\[ L = \sum_n \frac{-a}{2g^2} \text{Tr} \left( U^{-1} \dot{U} U^{-1} \right) - V \{ U \}, \]
where \( V \) contains the spacelike placquette terms, without any time derivatives.

To get the Hamiltonian we need first to find the canonical momenta \( \Pi = \frac{\delta L}{\delta \dot{q}}. \) In our case the coordinates \( q \) are the matrix elements of \( U \) on each spacelike link \( U_{\ell ab}, \) where \( \ell \) specifies the link (both spacelike direction and spatial position) and \( ab \) are the matrix indices. For convenience we will define the matrix \( P \) as the transpose, so
\[ \left( P_\ell \right)_{cb} := \frac{\delta L}{\delta (U_{\ell})_{bc}} = -\frac{a}{g^2} \left( U_{\ell}^{-1} \dot{U}_\ell U_{\ell}^{-1} \right)_{cb}, \]
\[ H = \sum_{\ell, c, b} \dot{U}_{\ell cb} \dot{P}_{\ell cb} - L = \sum_{\ell} \frac{-g^2}{2a} \text{Tr} \left( U_t P_t U_t P_t \right) + V \{ U \}. \]

Quantum mechanically, we write the momentum conjugate to \( q \) as \( -i \frac{\partial}{\partial q}, \) so \( P_{t cb} = -i \frac{\partial}{\partial q}. \) But \( P \) only enters \( H \) together with \( U. \) So let us define, for each link,
\[ \left( E_\ell \right)_{ab} = i \left( U_t P_t \right)_{ab} = U_{\ell ac} \frac{\partial}{\partial (U_{\ell})_{bc}}. \]
In terms of of the \( E \)’s, the kinetic energy term in \( L \) becomes \( \sum_\ell \frac{a}{2g^2} \text{Tr} E_\ell^2. \)

The \( E \)’s are a scaled electric field (\( ag \) times the usual field). As a differential operator it is very interesting.

But before we pursue the group theoretic arguments over the interpretation of \( E, \) let us meet an objection to the above procedure, that the \( U_{ab} \) are not independent real coordinates\(^2\). Instead, we should write \( U = e^{i \omega^i L_i} \) and treat the real, independent \( \omega^i \) as the coordinates. Then we get
\[ \Pi_i = \frac{\delta L}{\delta \dot{\omega}^i} = \frac{\delta U_{ab}}{\delta \dot{\omega}^i} \frac{\delta L}{\delta U_{ab}} = \frac{\delta U_{ab}}{\delta \omega^i} \frac{\delta L}{\delta U_{ab}}, \]
where I have made use of \( \dot{U}_{ab} = \frac{\delta U_{ab}}{\delta \omega^i} \dot{\omega}^i. \)

But
\[ \frac{\delta L}{\delta U_{ab}} = -\frac{a}{g^2} \left( U_{ab} \dot{U} \right), \]
so
\[ \Pi_i = -\frac{a}{g^2} \text{Tr} \left( \frac{\delta U_{ab}}{\delta \omega^i} U_{ab}^{-1} \left[ \sum_j \omega^j \frac{\delta U}{\delta \omega^j} U_{ab}^{-1} \right] U_{ab}^{-1} \right). \]
and
\[ \dot{\omega}^i \Pi_i = -\frac{a}{g^2} \text{Tr} \left( \dot{\omega}^i \frac{\delta U_{ab}}{\delta \omega^j} U_{ab}^{-1} \frac{\delta U}{\delta \omega^j} U_{ab}^{-1} U_{ab}^{-1} \right) = -\frac{a}{g^2} \text{Tr} \left( \dot{U} U_{ab}^{-1} U_{ab}^{-1} \right). \]
\(^2\)Of course we have a \( U \) and an \( \omega \) for each link, but we will drop the index \( \ell, \) or \( \bar{a}, \) that specifies which link, in what follows.
in agreement with what we found by careless manipulation.

One sees, however, that there are only $N^2 - 1$ coordinates and momenta for $SU(N)$, not $N^2$ as on might expect from $P_{\ell cb}$. The requirement that $U$ be unitary and of determinant 1 means there are constraints on the $P_{\ell cb}$’s — we can vary $U$ only preserving the unitarity and determinant. Interpreting the $P_{\ell cb}$ quantum mechanically as $-i\frac{\delta}{\delta U_{bc}}$ would be varying into a dimension for which the action is undefined, where $\det U \neq 1$. Our $E_{\ell cb}$ do not have that problem. Consider any function $f$ on $n \times n$ unitary matrices $u$. Each such matrix is $u = e^{i\phi}U$ with $U \in SU(N)$, so we may write $f(u) = f(\phi, U)$ and

$$\frac{\partial f}{\partial U_{ab}}|_{\phi} = \frac{\partial f}{\partial u_{ac}} \frac{\partial u_{ac}}{\partial U_{ab}}|_{\phi} = e^{i\phi} \frac{\partial f}{\partial u_{ab}},$$

so if $E_{ab} = iu_{ac} \frac{\partial}{\partial u_{bc}}$, $E_{ab} f = i e^{i\phi} U_{ac} \frac{\partial f}{\partial U_{bc}} = iU_{ac} \frac{\partial f}{\partial U_{bc}}|_{\phi}$ and has no dependence on how $f$ varies off the $\det u = 1$ subset. So we see that the $E_{ab}$ does not involve a derivative in the direction away from $\det U = 1$, but does form a complete set of $N^2 - 1$ first order differential operators on functions of $SU(N)$ matrices.

The $N^2 - 1$ momenta conjugate to $\omega^i$ are quantum mechanically

$$\Pi_i = -i \frac{\delta}{\delta \omega^i} = -i \frac{\delta U_{ab}}{\delta \omega^i} \frac{\delta}{\delta U_{ab}}.$$ Collectively these clearly span the same space as the $E_{ab}$. To understand the connection, consider different ways of asking how a function on $G$ varies as we move the element $g \in G$.

We see that $\Pi_i$ is not the same as $E_i$. To understand the connections, consider a function $f : G \rightarrow C$ defined on the group. Then derivatives of $f$ mean finding the change in $f(U)$ under small changes in $U$. If we write $U = e^{i\omega^L_i L_i}$, then

$$\frac{\partial}{\partial \omega^i} f(U) = i\Pi_i f.$$ On the other way, we could consider the change in $f$ as $U \rightarrow U'(\omega) = e^{i\omega^L_i L_i} U$, for infinitesimal $\omega^i$. If we define the differential operators

$$E_{ij} f(U) = \left. -i \frac{\partial}{\partial \omega^i} f \left( e^{i\omega^L_j L_j} U \right) \right|_{\omega = 0} = \left. -i \frac{\partial f(U')}{\partial U_{ab}} \frac{\partial U_{ab}}{\partial \omega_j} \right|_{\omega_j = 0} = \left. \frac{\partial f(U)}{\partial U_{ab}} i(L_{ij})_{ac} U_{cb} \right|_{\omega = 0} = (L_{ij})_{ac} U_{cb} \frac{\partial}{\partial U_{ab}} f = \text{Tr}(L_j E) f.$$  

So we see that the electric fields $E_{ab} = U_{\ell ac} P_{\ell cb}$ which we defined earlier are just linear combinations of the left-handed differential operators $E_j$, with $E = \sum_j E_j L_j$.

For completeness, let us also define a right derivative on the space of functions from $G \rightarrow C$ by

$$R_j f(U) = -i \frac{\partial}{\partial \omega^i} f \left( U e^{i\omega^L_j L_j} \right) \bigg|_{\omega = 0} = \frac{\partial f}{\partial U_{cb}} \left( -i \frac{\partial}{\partial \omega^j} \left( U e^{i\omega^L_j L_j} \right) \right)_{cb} \bigg|_{\omega = 0} = \frac{\partial f}{\partial U_{cb}} U_{ca} (L_j)_{ab}.$$ This is not exactly the same as $E_j$. If we make a matrix from $R_j$ in the same way as we did for $E_j$,

$$R_{ab} = \sum_j R_j \tau_{ja} \delta \frac{\delta}{\delta U_{ca}},$$

and recall $E_{ab} = U_{ac} \frac{\delta}{\delta U_{cb}}$, so

$$R f = \delta \frac{\delta f}{\delta U} U = U^{-1} U \frac{\delta f}{\delta U} U = U^{-1} (Ef) U.$$ Note that the $E_j$, $R_j$ and $\Pi_j$ are each sets of $N^2 - 1$ first order differential operators on the space of functions from $G \rightarrow C$, an $N^2 - 1$ dimensional space. Thus they are linear combinations of each other. The connection between them is associated with the adjoint representation of the group.

We will write $[\mathcal{S}(L_l)]_{jk}$ for the adjoint representation of the generators, so $[\mathcal{S}(L_l)]_{jk} = \delta_{j}^{k}$. Let us call the adjoint representation of the group $\mathcal{S}$:

$$S_{jk} (e^{i\omega^L l}) = \left[ e^{i\omega^L \mathcal{S}(L_l)} \right]_{jk}.$$ [Note: $\mathcal{S}$ and $S$ were both previously called $\Gamma^{\text{adj}}$. Also it is hard to distinguish $\mathcal{S}$ from $S$. ] These will come up in the connections we are seeking.

### 14.1 Relation of $R_j$ to $E_j$

Now $U^{-1} L_l U$ is a linear combination of the $L$’s, so $U^{-1} L_l U = M_{lm}(U) L_m$. We can show $M_{tm} = S_{tm}$.

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Proof: Let $U(\lambda) = e^{i\lambda \omega^1 L_1}$, $M(\lambda) = M(U(\lambda))$. Then $e^{-i\lambda \omega^1 L_1} L_1 e^{i\lambda \omega^1 L_1} = M_{\ell m}(\lambda) L_m$. Differentiate with respect to $\lambda$:

$$e^{-i\lambda \omega^1 L_1} \left[ L_1, i \omega^k L_k \right] e^{i\lambda \omega^1 L_1} = \frac{dM_{\ell m}(\lambda)}{d\lambda} L_m = -\omega^k c_{\ell k n} M_{nm}(\lambda) L_m.$$ 

The $L_m$ are linearly independent, so

$$\left( \frac{dM(\lambda)}{d\lambda} \right)_{\ell m} = -\omega^k c_{\ell k n} M_{nm}(\lambda) = i\omega^k \left[ S(L_k) M \right]_{\ell m}.$$ 

This is a differential equation with solution

$$M(\lambda) = e^{i\lambda \omega^k S(L_k) K}$$

with initial condition $M_{\ell m}(0) = \delta_{\ell m} = \mathbb{I} = K$, so $M(\lambda) = e^{i\lambda \omega^k S(L_k)} = S(U(\lambda))$. Thus

$$e^{-i\lambda \omega^k L_k} L_1 e^{i\lambda \omega^k L_k} = S_{\ell m}(e^{i\lambda \omega^k L_k}) L_m = \left( e^{i\lambda \omega^k S(L_k)} \right)_{\ell m} L_m.$$ 

So $U^{-1} L_1 U = S_{\ell m}(U) L_m$ and by exponentiation, $U^{-1} e^{i\lambda \omega^k L_k} U = e^{i\lambda \omega^k S_{\ell m}(U)} L_m$. Now reconsider

$$R_\ell f(U) = -i \frac{\partial}{\partial \omega^\ell} f \left( U e^{i\lambda \omega^k L_k} \right) \bigg|_{\omega = 0}$$

$$= -i \frac{\partial}{\partial \omega^\ell} f \left( e^{i\lambda \omega^k S_{\ell m}(U^{-1}) L_m} U \right) \bigg|_{\omega = 0}$$

$$= -i \frac{\partial}{\partial \omega^\ell} \frac{\partial}{\partial \omega^m} f \left( e^{i\lambda \omega^m L_m} U \right) \bigg|_{\omega = 0}$$

where $\omega^m = \omega^k S_{\ell m}(U^{-1})$. So

$$R_\ell f(U) = S_{\ell m}(U^{-1}) E_m f(U), \quad R_\ell = S_{\ell m}(U^{-1}) E_m, \quad \text{and} \quad E_m = S_{\ell m}(U) R_\ell,$$

as $S^{-1}(U^{-1}) = S(U).$
or
\[
\frac{\partial m}{\partial \omega^k} = \int_0^1 d\alpha \left[ e^{-i\omega^k S(L_k)} \right]_{\ell m} = \left[ \frac{1 - e^{-i\omega^k S(L_k)}}{i\omega^k S(L_k)} \right]_{\ell m}.
\]
so
\[
\frac{\partial}{\partial \omega^k} = \left[ \frac{1 - e^{-i\omega^k S(L_k)}}{\omega^k S(L_k)} \right]_{\ell m} E_m
\]
gives the connection of the left derivative with the ordinary partial derivative. [Note: the fraction in square brackets is well defined even though $\omega^k S(L_k)$ is likely not an invertible matrix, so dividing by it is dubious. But the expression in the bracket has a well defined power series expansion and thus a well defined meaning.]

We have given a great deal of discussion to our Hamiltonian operator. We now turn to the states on which this operator acts. A state can be specified (a) at each site, the occupation number for the fermions at that site (b) on each link, a group element. We focus on this.

We are not accustomed to having dynamics with canonical coordinates defined on a compact space like $U \in G$, except in advanced courses in classical mechanics (e.g. the orientation of a rigid body). In quantum mechanics we know that the momentum is related to translations, so the ordinary kinetic energy term $\sum_i p_i^2/2m$ is connected with a space of $q$'s which is translation invariant in each $p_i$. Our kinetic energy is not of that form, as the $E_i$'s don't commute and the $\partial/\partial \omega^k$'s do not enter $H$ in this simple fashion. This is consistent with the fact that our space of $U$'s is just not a Euclidean $N^2 - 1$ dimensional space. Now quantum mechanics can be viewed as a functional integral over all possible configurations, which requires a weight for the configurations. In a deep sense this weight is connected with the kinetic energy term in the Lagrangian density. We can sidestep this issue (temporarily) by observing that the only sensible way to integrate over the $U$'s is in a group invariant way. In terms of eigenstates $|U_{ab}\rangle$ of $U$ (on a given link), we need to be able to write

\[
I = \int d\mu(U) |U_{ab}\rangle \langle U_{ab}|,
\]
or in terms of the parameters $\omega^k$, $d\mu(U) = d\mu(\omega^k)$, the measure must be group invariant, $\int d\mu(\omega)f(U(\omega)) = \int d\mu(\omega)f(GU(\omega))$ for any function $f : G \to \mathbb{C}$ and any fixed element $G$ in the group. This is also what we needed long ago to define the orthogonality of group representations.

So now we leave physical applications and turn to the problem of finding the Haar or Hurwitz measure $d\mu(\omega^k)$ satisfying the invariance.

Let us write $d\mu(\omega) = h(\omega)d\omega^1 \cdots d\omega^n$ for an $n$ dimensional Lie group. For a fixed $G$ and $\omega$, the group element $U = GU(\omega) = U(\omega') = e^{i\omega' L_j}$, where the $\omega'$'s are functions of the $\omega$'s (and also of $G$, but we are holding $G$ fixed).

If we change variables on the right hand side of our invariance equation, we get

\[
\int h(\omega)f(U(\omega)) d\omega^1 \omega' = \int h(\omega)f(U(\omega')) \frac{\partial \omega'^j}{\partial \omega^k} d\omega^1 \omega',
\]
where $|\partial/\partial| is the Jacobian determinant. Invariance for arbitrary functions $f$ clearly requires

\[
h(\omega') = h(\omega) \left. \frac{\partial \omega'^j}{\partial \omega^k} \right|_{\omega'=\omega}.
\]
Let $h(0) = c$. Then $h$ at any other point $\omega$ is determined by taking $G = e^{i\omega L_j}$,

\[
h(\omega) = c \left. \frac{\partial \omega'^k}{\partial \omega^j} \right|_{\omega'=0}^{-1},
\]
where $e^{i\omega L_k} = e^{i\omega' k} e^{i\omega L_k}$.

This is a necessary condition for the invariance. To show it is also sufficient, we need to show that the relation holds for all $\omega$. Consider the measure at two points, $\omega'$ and $\omega''$. We need to show that if $U(\omega'') = GU(\omega')$ for a fixed $G$,

\[
h(\omega'') = h(\omega') \left. \frac{\partial \omega''}{\partial \omega'} \right|_{\omega''=\omega'}.
\]
Let $G' = e^{i\omega' L_k}$ and $G'' = e^{i\omega'' L_k}$ and $U(\nu') = G'U(\nu)$ and $U(\nu'') = G''U(\nu)$. Then $U(\nu'') = G''(G')^{-1}U(\nu') = GU(\nu')$, so $G = G''(G'^{-1})$,

\[
\left. \frac{\partial \omega''}{\partial \omega'} \right|_{\omega''=\omega'} = \left. \frac{\partial \omega'}{\partial \omega'} \right|_{\nu'} / \left. \frac{\partial \nu'}{\partial \omega'} \right|_0 = \left( \frac{c}{h(\omega'')} \right) / \left( \frac{c}{h(\omega'')} \right) = \frac{h(\omega')}{h(\omega'')},
\]
which verifies the invariance in general.
Now we need to evaluate the Jacobian at \( \nu = 0 \). As \( e^{i\nu^k L_k} = e^{i\omega^k L_k} e^{i\nu^k L_k} \),
\[
\frac{\partial}{\partial \nu^k} e^{i\nu^k L_k} \bigg|_{\nu=0} = \frac{\partial}{\partial \nu^j} e^{i\nu^k L_k} \frac{\partial}{\partial \nu^j} e^{i\nu^k L_k} \bigg|_{\nu=0} = \frac{\partial \nu_j}{\partial \nu^k} e^{i\nu^k L_k} iL_j
\]
so
\[
\frac{\partial \nu_j}{\partial \nu^k} L_j = \int_0^1 d\alpha e^{i(1-\alpha)\nu^k L_k iL_j} e^{i\alpha \nu^k L_k} \bigg|_{\nu'=\omega} = \int_0^1 d\alpha e^{-i\alpha \omega^k L_k} L_k e^{i\alpha \nu^k L_k} = \left[ \frac{1-e^{i\omega^m S(L_m)}}{i\omega^m S(L_m)} \right]_{kj} L_j.
\]
Thus
\[
h(\omega) = c \det \frac{\partial \nu}{\partial \nu'} = c \det \left( \frac{e^{i\omega^m S(L_m)} - 1}{i\omega^m S(L_m)} \right).
\]
For any fixed (real) \( \omega \), \( \omega^m S(L_m) \) is a hermitean matrix and can be diagonalized. If \( \lambda_j \) are the real eigenvalues, we have
\[
h(\omega) = c \prod_j \frac{e^{i\lambda_j} - 1}{i\lambda_j} = \prod_j \frac{\sin(\lambda_j/2)}{\lambda_j/2} \prod_j e^{i\lambda_j/2}.
\]
But \( \omega^m S(L_m) \) is also imaginary, while still hermitean, so it is antisymmetric and has zero trace. This is preserved by similarity, so \( \sum_j \lambda_j = 0 \), and
\[
h(\omega) = c \prod_j \frac{\sin(\lambda_j/2)}{\lambda_j/2}.
\]

Let us consider an explicit example, \( SU(2) \). Any \( U = e^{i\omega^z z/2} \) is equivalent (by similarity) to one along the \( z \) axis, with the same \( \omega^z \). Then \( \omega^z S(L_z) = |\omega| S(L_3) \). \( S \) is the adjoint representation, so \( L_3 \) has the eigenvalues \( \pm 1, 0 \). The zero eigenvalue gives 1 in the product (recall the note on the definition of the \( [\ ] \) expression) and the others each give 2 \( \sin(|\omega|/2)/|\omega| \), so all together
\[
h(\omega) = 4c \frac{\sin^2(|\omega|/2)}{\omega^2}.
\]
Choosing \( c = 1 \), we may take the invariant volume to be
\[
\int_{|\omega| \leq 2\pi} d\omega^1 d\omega^2 d\omega^3 \frac{4 \sin^2(|\omega|/2)}{|\omega|^2}.
\]

If we used spherical coordinates for \( \omega \), this gives
\[
\int_0^{2\pi} d\omega^1 d\omega^2 \omega^3 \frac{4 \sin^2(|\omega|/2)}{|\omega|^2}.
\]
where \( d\Omega = \sin \theta d\theta d\phi \) as usual. We note that the volume with \( |\omega| \geq 2\pi - \epsilon \) is \( \frac{4\pi^3}{3} e^3 \), the same as the volume with \( |\omega| \leq \epsilon \), and not \( \epsilon \) times the ordinary area of the sphere, that is, \( 4\pi(2\pi)^2\epsilon \). The metric says points with \( \omega \approx 2\pi \), even in very different directions, are close together, because for \( \omega = 2\pi \), all these \( \omega \)'s correspond to the same \( U = \begin{pmatrix} -1 & 0 \\ 0 & -1 \end{pmatrix} \). Note \( h(\omega) \) is exactly the metric\(^3\) on \( S^3 \) induced by embedding in \( E^4 \), with the angle with respect to the 4th axis \( \omega/2 \).

Some comments:

(a) For \( SU(3) \), not all directions are equivalent. So there is no equivalent of \( |\omega| \).

(b) For many applications, we needed the existence of the invariant measure but not the actual expression for it. For example, in our orthogonality proofs.

(c) Most discussions of \( SU(2) \) parameterize the group in terms of Euler angles rather than in the way we did, the angle of the rotation and the direction of the axis.

(d) Note that the factor in \( h \) is just the determinant of the matrix which connects the left derivatives to the ordinary derivatives. This is not an accident, but rather has to do with this relation defining a metric (not just a measure) on the group space, and the measure given by the determinant of the metric.

\(^3\)See “hypersph” in Supplementary Notes.