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Representation of Group Generators by Boson or Fermion Operators Application to Spin Perturbation Theory

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(27. V. 66)

A representation of spin by fermion operators was known for a long time for $S = 1/2$ [1]¹), whereas the generalization to $S > 1/2$ is more recent [2]. The interest of such a representation is to avoid the complications of a generalized Wick's theorem valid for spin [3]. Here we remark first that such representations have a much larger generality, valid for fermions as well as for bosons [4]. We then discuss the elements of perturbation theory in the representation of spin by fermion operators.

1. Let G_i ($i = 1, \dots, r$) be the generators of a simple Lie group numbered such that the first l commute with each other. A basis of an irreducible representation D^N may then be chosen such that

$$G_i |\nu\rangle = m_i(\nu) |\nu\rangle; \quad i = 1, \dots, l; \quad \nu = 0, 1, \dots, N-1. \quad (1)$$

One calls r the order and l the rank of the group, N the dimension and $m_i(\nu)$ the weights of the representation [5].

Let a_ν and a_ν^* be boson or fermion operators such that

$$[a_\nu a_{\nu'}^*]_\mp = \delta_{\nu \nu'}. \quad (2)$$

It is then easy to prove that the expressions $\sum_{\nu \nu'} a_\nu^* g_i^{\nu \nu'} a_{\nu'}^* (i = 1, \dots, r)$ satisfy the commutation rules of the G_i , $[G_i G_j]_- = \sum_k \gamma_{ij}^k G_k$, if the matrices g_i do so.

If $| \rangle$ is the vacuum state of these fictitious particles such that

$$a_\nu | \rangle = 0 \quad (3)$$

then the identification

$$|\nu\rangle = a_\nu^* | \rangle \quad (4)$$

$$G_i = \sum_{\nu \nu'} a_\nu^* g_i^{\nu \nu'} a_{\nu'}^* \quad (5)$$

reproduces all the properties of the representation (1), of which the matrices g_i are a realization.

The inconvenience of this representation by *fictitious* particles ν is that neither the vacuum $| \rangle$ nor the states with more than one particle, $a_{\nu_1}^* \dots a_{\nu_n}^* | \rangle$, $n > 1$, have

¹⁾ Numbers in brackets refer to References, page 465.

physical meaning. For fermions $n \leq N$ and hence the dimension of the Hilbert space of all particles is 2^N .

The case of spin is obtained with the group $SU(2)$ for which $r = 3$, $l = 1$, and $G_1 = S_3$, $G_2 = S_+$, $G_3 = S_-$, $N = 2S + 1$, $m(\nu) = S - \nu$.

For spins localized on different atoms labeled by n one has [2]

$$S_n = \sum_{\nu \nu'} a_{\nu n}^* s^{\nu \nu'} a_{\nu' n}. \quad (6)$$

It is interesting to note that in elementary particle theory, localized G_i have an interpretation as generalized charge densities and (5) expresses their bilinear form in the fields of the particles ν which, in this case, are *physical*.

2. Consider now perturbation theory for the coupling of one single spin, Equation (6) with indices n omitted. The unperturbed hamiltonian is, for an external magnetic field \mathfrak{H} ,

$$H_0 = -\gamma S_z = -\gamma \sum_{\nu=0}^{2S} (S - \nu) a_{\nu}^* a_{\nu}; \quad \gamma = 2\mu_B \mathfrak{H} > 0. \quad (7)$$

$|0\rangle = a_0^* |0\rangle$ is the ground state. The expressions to be calculated are of the form $\langle \prod_{\nu=0}^{2S} \tau_{\nu} \rangle$. Here τ_{ν} is a product ordered according to imaginary times of operators $a_{\nu}^*(-i\tau)$, $a_{\nu}(-i\tau)$ where $O(t) = e^{iH_0 t} O e^{-iH_0 t}$. $\langle \rangle$ is an unperturbed canonical average taken over the physical states $|\nu\rangle$, i.e.

$$\langle O \rangle \equiv \sum_{\nu=0}^{2S} \langle \nu | e^{-\beta H_0} O | \nu \rangle / \sum_{\nu=0}^{2S} \langle \nu | e^{-\beta H_0} | \nu \rangle. \quad (8)$$

One calculates

$$\langle \nu | e^{-\beta H_0} \prod_{\nu'} \tau_{\nu'} | \nu \rangle = e^{\beta \gamma (S - \nu)} \langle \nu | \tau_{\nu} | \nu \rangle \prod_{\nu' \neq \nu} \langle | \tau_{\nu'} | \rangle. \quad (9)$$

Here $\langle | \tau_{\nu} | \rangle$ is obtained by applying the usual diagram technique for *zero temperature* and is expressed in terms of free propagators which, however, do not have the usual form because of the imaginary time order. They are

$$\begin{aligned} G_{\nu}(\tau) &\equiv \langle | T_{\tau} (c_{\nu}(-i\tau) a_{\nu}^*(0)) | \rangle \\ &= \frac{1}{2\pi} \int_{-i\gamma(S-\nu+0)-\infty}^{-i\gamma(S-\nu+0)+\infty} d\omega \frac{e^{i\omega\tau}}{i\omega - \gamma(S-\nu)} = \begin{cases} e^{\tau \gamma (S - \nu)}; & \tau > 0 \\ 0; & \tau < 0 \end{cases}. \end{aligned} \quad (10)$$

The expression $\langle \nu | \tau_{\nu} | \nu \rangle$ is not directly accessible to the usual technique because the normal products obtained by applying the ordinary Wick's theorem to τ_{ν} , do not all vanish when taken between the states $|\nu\rangle$. To handle it, we introduce the unperturbed canonical average taken over the Hilbert space of the fictitious particles. For fermions this is

$$\begin{aligned} \langle\langle O \rangle\rangle &\equiv \sum_{\{n_{\nu}=0,1\}} \langle | \prod_{\nu} (a_{\nu})^{n_{\nu}} e^{-\beta H_0} O \prod_{\nu} (a_{\nu}^*)^{n_{\nu}} | \rangle \\ &\times \left\{ \sum_{\{n_{\nu}=0,1\}} \langle | \prod_{\nu} (a_{\nu})^{n_{\nu}} e^{-\beta H_0} \prod_{\nu} (a_{\nu}^*)^{n_{\nu}} | \rangle \right\}^{-1}. \end{aligned} \quad (11)$$

One then finds that $\langle \nu | \tau_\nu | \nu \rangle$ can be expressed in terms of $\langle | \tau_\nu | \rangle$ and $\ll \tau_\nu \gg$,

$$e^{\beta \gamma (S-\nu)} \langle \nu | \tau_\nu | \nu \rangle = (1 + e^{\beta \gamma (S-\nu)}) \ll \tau_\nu \gg - \langle | \tau_\nu | \rangle. \quad (12)$$

$\ll \tau_\nu \gg$ can be calculated by the usual diagram technique for *finite temperature* [6] and is expressed in terms of free propagators,

$$g_\nu(\tau) \equiv \ll T_\tau (a_\nu(-i\tau) a_\nu^*(0)) \gg = \frac{1}{\beta} \sum_{r=-\infty}^{+\infty} \frac{e^{i\omega_r \tau}}{i\omega_r - \gamma(S-\nu)}; \quad \omega_r = \frac{\pi}{\beta} (2r+1). \quad (13)$$

The general result now is

$$\begin{aligned} \langle \prod_\nu \tau_\nu \rangle &= e^{-\beta \gamma S} \frac{1 - e^{-\beta \gamma}}{1 - e^{-\beta \gamma (2S+1)}} \left\{ \sum_{\nu=0}^{2S} (1 + e^{\beta \gamma (S-\nu)}) \right. \\ &\quad \times \ll \tau_\nu \gg \prod_{\nu' \neq \nu} \langle | \tau_{\nu'} | \rangle - (2S+1) \langle | \prod_\nu \tau_\nu | \rangle \left. \right\}. \end{aligned} \quad (14)$$

Since $\prod_\nu \tau_\nu$ usually is a product of operators (6) one has $\langle | \prod_\nu \tau_\nu | \rangle = 0$. For low temperatures, $\beta \gamma \gg 1$, Equation (14) simplifies considerably:

$$\langle \prod_\nu \tau_\nu \rangle \simeq \ll \tau_0 \gg \prod_{\nu \neq 0} \langle | \tau_\nu | \rangle. \quad (15)$$

The procedure described here is to be compared with that of Abrikosov's second paper [7] while in his first paper [2], ABRIKOSOV uses an artificial hamiltonian which is obtained from (7) by the substitution $\gamma(S-\nu) \rightarrow \lambda$. In this case, Equation (14) goes over into

$$\langle \prod_\nu \tau_\nu \rangle = \frac{e^{\beta \lambda} + 1}{2S+1} \sum_\nu \ll \tau_\nu \gg \prod_{\nu' \neq \nu} \langle | \tau_{\nu'} | \rangle. \quad (16)$$

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$$|\nu\rangle = [(2S-\nu)! \nu!]^{-1/2} (a_0^*)^{2S-\nu} (a_1^*)^\nu |0\rangle$$
- instead of Equation (4).
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