mein-physikalischer Wichtigkeit sein kann, als die Bedingungen bei der Ableitung der Maxwellschen Gleichungen aus den tieferliegenden Gleichungen des metrischen Feldes eingehalten sind. Diese Bedingungen waren:

1. Die Zeit ist kontinuierbar, was heißt, daß Prozesse, die schneller als in etwa $10^{-20}$ sec ablauen, nicht erfaßt werden.


On the Quantum Theory of the Anharmonic Oscillator in Functional Space

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(Z. Naturforsch. 24 a, 188—197 [1969] ; received 17 October 1968)

Dynamics of quantum field theory can be formulated by functional equations. For strong interaction nonperturbative solutions of these functional equations are required. For the investigation of solution procedures the model of an anharmonic oscillator is used, because of its structural equivalence with dressed one- and two-particle states of field theory. To perform a variational solution procedure a scalar product for the state functionals is introduced and its existence is proven. The scalar product definition admits a mapping of the physical Hilbert space on the functional space. Therefore a "functional" quantum theory seems to be possible. The whole procedure can be transferred to relativistic invariant field theories, provided these theories are regularized to give finite results at all.

In quantum field theory especially in nonlinear spinor theory of elementary particles the dynamical behaviour of the physical systems can be described by functionals of field operators in a Heisenberg representation and corresponding functional equations. To obtain the physical information in the strong coupling case, it is necessary to solve the functional equations without perturbation theory. As has been discussed in previous papers the anharmonic oscillator is a suitable test problem for the investigation of solution procedures of strong coupling functional equations. Therefore in the following this model will be treated. The general idea for the nonperturbative solution of the functional equations is the use of an expansion of the physical functionals into series of suitably chosen base functionals and to approximate the exact infinite series by series with a finite number of terms. The complete theory then requires the proof of convergence and the explicit calculation of the approximate functionals. For the limiting case of single time functionals, being ordinary functions in the oscillator model both problems have been discussed. For the general time ordered functionals only the second problem has been investigated in preceding papers. In this paper now we attack the convergence problem. As is well known from functional analysis a necessary condition for convergence is the
existence of a scalar product definition. Therefore the main effort is made in this direction. Guided by the idea, that the entire quantum theory should be expressible in a selfconsistent way as a functional quantum theory, we propose a certain scalar product definition and prove its existence for the exact physical state functionals. As a consequence a mapping of the physical Hilbert space on the functional Hilbert space is possible and one is able to formulate quantum theory in functional space only. Because the scalar product definition proposed is explicitly a many-time formalism one may use it for canonical as well as for noncanonical quantized theories. It is especially this feature which is of interest for nonlinear spinor theory of elementary particles, because so far no proper scalar product definition has been given there. The treatment of the problem stated here is by no means complete, so further papers with theoretical as well as numerical analysis are in preparation.

1. Fundamentals

Identifying \( q(t) \) with \( \psi_1(t) \) and \( p(t) \) with \( \psi_2(t) \) the equations of motion for an anharmonic oscillator can be written according to (1) in the general form

\[
\frac{d}{dt} \psi_\alpha(t) = B_{\alpha\beta} \psi_\beta(t) + C_{\alpha\beta} D_\gamma \psi_\beta(t) \psi_\gamma(t) \psi_\epsilon(t)
\]

with the commutation relation

\[
[\psi_\alpha(t), \psi_\beta(t)] = i A_{\alpha\beta} \mathbf{1}.
\]

This representation is in complete analogy to the nonlinear spinor equation of elementary particle theory with Hermitean field operators. The matrices \( B, C, D \) are to a certain extent still arbitrary, while \( A \) is fixed by the canonical quantization to be

\[
A_{\alpha\beta} := \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}.
\]

A very suitable anharmonic oscillator model is given by the following choice \( B = C = A \) and \( D = \delta_{\gamma\alpha} \).

Its stationary energy eigenstates are

\[
|n\rangle := f_n(\psi_1)
\]

with \( \psi_1 \equiv \psi_1(0) \) and

\[
f_n(x) = \pi^{-1} (2\pi n!)^{-1} e^{-\frac{1}{2} x^2} H_n(x)
\]

where \( H_n \) are the standard Hermitean polynomials. The corresponding energies are

\[
E_n = (n + \frac{1}{2}) + (n + \frac{1}{2})^2.
\]

For the field theoretic functional treatment we characterize the stationary states of the model by state functionals

\[
\mathcal{Z}_a(j) := \langle 0 | T \exp \int \psi_a(\xi) j_a(\xi) d\xi | a \rangle
\]

where \( j_a(\xi) \) are commuting functional source-functions. Using the procedure outlined in II, App. I, we obtain for the state functionals the functional equations

\[
\frac{d}{dt} \mathcal{Z}_a(j) = \left[ C_{\alpha\beta} D_\gamma i \frac{\delta}{\delta j_\alpha(t)} i \frac{\delta}{\delta j_\beta(t)} + B_{\alpha\beta} i \frac{\delta}{\delta j_\beta(t)} - A_{\alpha\beta} j_\beta(t) \right] \mathcal{Z}_a(j)
\]

and the subsidiary condition

\[
P \mathcal{Z}_a(j) := \int j_a(t) \frac{d}{dt} i \frac{\delta}{\delta j_a(t)} dt \mathcal{Z}_a(j) = -\omega_a \mathcal{Z}_a(j)
\]

with \( \omega_a = (E_a - E_0) \). Combining Eqs. (1.8) and (1.9) we obtain the eigenvalue equation

\[
\int j_a(t) \left[ C_{\alpha\beta} D_\gamma i \frac{\delta}{\delta j_\alpha(t)} i \frac{\delta}{\delta j_\beta(t)} + B_{\alpha\beta} i \frac{\delta}{\delta j_\beta(t)} - A_{\alpha\beta} j_\beta(t) \right] dt \mathcal{Z}_a(j) = -\omega_a \mathcal{Z}_a(j)
\]

where the last term on the right side of (1.8) drops out. For simplicity we shall work in the following with Eq. (1.10). In nonlinear spinor theory it turns out, that this equation is not explicitly relativistically invariant and therefore one has to use (1.8) and (1.9). But all arguments applied to (1.10) are valid as well as for (1.8) and (1.9), so no restriction is imposed on the entire problem by using (1.10). For the further investigation we have to expand the functionals (1.7) into power series. We obtain

\[
\mathcal{Z}_a(j) = \sum_{k=1}^{\infty} \frac{i^k}{k!} \int \tau_k(t_1 \ldots t_k) j_\alpha(t_1) \ldots j_\alpha(t_k) dt_1 \ldots dt_k
\]

\[\]
which follows immediately from the exponential series of (1.7).

2. Rigged Hilbert Space Representation

In the preceding section the state functionals were introduced only formally. For working with them we have to achieve their real existence. Thereby it is more convenient to define the state functionals by the expansion (1.11). Because in (1.11) the \( \tau_k \) functions are distributions we have to restrict the source functions \( j_\alpha (\xi) \) in a suitable way in order to obtain at least finite values of the integrals. A special choice for the source functions is the space \( S \) of all rapidly decreasing and infinitely often differentiable functions. Introducing a complete set of orthonormalized base functions in this space, one may expand the source functions within this set. Without any loss of generality we may choose the oscillator functions (1.5) to be this set. Then we have

\[
\tau_\alpha (t_1 \ldots t_n) = \sum_k \phi_{\alpha k} f_k (t) \tag{2.1}
\]

and by substituting (2.1) into (1.11) the functionals are given by

\[
\Xi (j) = T (\phi) = \sum_n \sum_{t_1 \ldots t_n} \frac{i^n}{n!} T \left( \tau_1 \ldots \tau_n \right) \phi_{1 t_1} \ldots \phi_{n t_n} \tag{2.2}
\]

with

\[
T \left( \tau_1 \ldots \tau_n \right) = \int \tau_\alpha (t_1 \ldots t_n) f_{\alpha} (t_1) \ldots f_{\alpha} (t_n) \, dt_1 \ldots dt_n \tag{2.3}
\]

In App. II the existence and a numerical estimate of (2.3) is proven. By (2.3) one learns, that this representation means a spectral decomposition of the set (1.12) by \( L_2 \) functions. Because the functions (1.12) are distributions this spectral decomposition leads to a rigged Hilbert space representation of the \( \tau_\alpha \) functions. Using the completeness relation of the \( f_\alpha \) functions from (2.3) they follow to be

\[
\tau_\alpha (t_1 \ldots t_n) = \sum_{u_1 \ldots u_n} T \left( \tau_1 \ldots \tau_n \right) f_{\alpha_1} (t_1) \ldots f_{\alpha_n} (t_n) \tag{2.4}
\]

This representation can be proven also by direct evaluation of the matrix representation of the \( \tau_\alpha \) functions.

For working with such a representation one has to transform also the eigenvalue equation resp. the other dynamical equations. By the completeness relations follows

\[
\frac{\partial}{\partial \tau_\alpha (t)} = \sum_{w_1 \ldots w_n} \sum_{k \neq w} \left( \frac{\partial}{\partial \phi_{w k}} \right) f_{w} (t) \tag{2.5}
\]

and the transformed eigenvalue equation (1.10) reads

\[
\left[ \sum_{k_1 \ldots k_4} \sum_{k_5 \ldots k_8} \phi_{k_1 k_2} \phi_{k_3 k_4} \phi_{k_5 k_6} \phi_{k_7 k_8} F \left( \tau_1 \ldots \tau_4 \right) \right. \\
\left. + B_{\alpha \beta} \sum_k \phi_{\alpha k} \phi_{\beta k} \right] T_\alpha (\phi) = 0 \tag{2.6}
\]

with

\[
F \left( \tau_1 \ldots \tau_4 \right) = C_{\tau_1 \tau_2} D_{\tau_3 \tau_4} \int f_{\alpha} (t) \, dt \tag{2.7}
\]

Additionally the state functionals have to satisfy the condition of stationarity (1.9), which reads in the rigged Hilbert space representation

\[
P T_\alpha (\phi) = -\omega_\alpha T_\alpha (\phi) \tag{2.8}
\]

with

\[
P := \sum_{k \neq w} \left[ \frac{\sqrt{k+1}}{2} \phi_{k+1, w} + \frac{\sqrt{k}}{2} \phi_{k-1, w} \right] \frac{3}{\sum_{k \neq w} \left[ \phi_{w k} \phi_{w \beta} \right]} \tag{2.9}
\]

In nonlinear spinor theory the decomposition (2.1) corresponds to a Lorentz invariant representation of \( f_\alpha (x) \) according to the irreducible representations of the rotation group. These representations are available. Therefore the rigged Hilbert space concept has a meaning in field theory too and all calculations made for the anharmonic oscillator in this representation can be applied equally well in field theory.
3. Norm Definition and Existence

Formation of a functional norm requires functional integration. To avoid all difficulties connected with the definition of a continuous functional integration, we perform functional integration in the rigged Hilbert space representation. Then this integration runs over a countably infinite set of variables and can be evaluated by algebraic methods. This is performed in App. I. Naturally the results of a genuine functional integration should be the same, but we do not discuss this problem here. Because the functionals (2.2) are not simply square summable for a proper norm definition weighting factors have to be introduced. We shall use two such factors. The first factor is derived functionally by the observation that for the operator

$$s_{ij} (3-1)$$

the functional equation

$$N a_i . . . Q \ln a_n = n q_{i1} z_i . . . q_{in_n}$$

holds, i.e. the power functions of $q_{in}$ with arbitrary arguments are eigenfunctions of this operator. Also any steady function of $N$ is defined in application to this eigenvector system, and can be applied therefore to $T(q)$ too. Using this property we define the first weighting factor by

$$r \equiv (V) T(e) (3.3)$$

where $dr(q)$ means simply Riemannian integration over all $q$-variables. Again the numerical values of $G_{na mb}$ are still open. In the next section we try to fix these values by physical considerations. In this section we prove only under very general assumptions about $G_{na mb}$ the existence of (3.4). To avoid unnecessary complication of the formulas, we introduce superindices, i.e. we denote the spinorial indices $z$ and the expansion indices $k$ of (2.1) by only one superindex. Then the norm formula (3.4) reads explicitly

$$\| T(q) \|^2 = \sum_{k, l} T^k (s_1 . . . s_k) (\frac{-i}{k!}) \langle D_k (s_1 . . . s_k, G) D_l (t_1 . . . t_l, G) \rangle \frac{i^l}{(l!)} T_l (t_1 . . . t_l)$$

(3.5)

where the Dyson functionals $^8 D_k$ in superindices are defined by

$$D_k (h_1 . . . h_k, G) := \frac{1}{k!} q_h \ldots \sum_{q_{n}} \exp \{- \frac{1}{2} \sum_{q_{n}} G_{n n m} q_m \}$$

(3.6)

and the brackets symbolize functional integration. In App. I the following formula for the functional integral of $D_k$ and $D_l$ is derived

$$\langle D_k (s_1 . . . s_k, G) D_l (t_1 . . . t_l, G) \rangle = \det U_{\min(k, l)} \sum_{r=0}^{\min(k, l)} C^{k+l}_{k-l} \sum_{s_1 . . . s_k} R_{s_1 s_2 . . . s_k} \sum_{t_1 . . . t_l} R_{t_1 t_2 . . . t_l} (3.7)$$

where $R$ is equal to $(\mathcal{R} G)^{-1}$ and the matrix $U$ is connected with the spectral representation of $\mathcal{R} G$. Details are given in App. I. From App. I follows, that $\mathcal{R} G$ has to have a spectral representation with all eigenvalues unequal zero. Then $(\mathcal{R} G)^{-1}$ exists also.

Substituting (3.7) into (3.5) and observing the permutational symmetries of the $T_k$-expansion coefficients (3.5) becomes

$$\| T(q) \|^2 = \det U \sum_{k, l} \sum_{s_1 . . . s_k} \sum_{t_1 . . . t_l} (-1)^k (k! \ l!)^r T^k (s_1 . . . s_l) T_l (t_1 . . . t_l) R_{s_1 t_1} R_{s_2 t_2} . . . R_{s_k t_k} (3.8)$$

For the further evaluation of (3.8) we have to make an estimate of $T_k$. This is done in App. II. One obtains

$$| T_k (s_1 . . . s_k) | \leq n^{4k} (k!)^{3/4} C_n$$

(3.9)

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By this approximation (3.8) goes over into the inequality
\[ \| T(q) \|^2 \leq \sum_{k,l} \sum_{r=0}^{\min(k,l)} C_r^{kl} \frac{(k!)^{2r}}{(k!l!)^{r-r/2}} K^{(k+l)} \| \det U \| \]
with
\[ K := \left( \sum_{u,w} R_{uw} \right)^2 C \]
and according to App. I
\[ C_r^{kl} = \frac{1}{\binom{k-r}{2} \binom{l-r}{2} r!} (\frac{1}{2})^{k+l-r} \]
Therefore
\[ \| T(q) \|^2 \leq \sum_{k,l} \sum_{r=0}^{\min(k,l)} \frac{K^{(k+l)}(k!)^{2r}}{(k!l!)^{r-r/2}} \frac{1}{\binom{k-r}{2} \binom{l-r}{2} r!} (\frac{1}{2})^{k+l-r} \| \det U \| \]
is obviously convergent as long as (3.11) is finite, \( \det U \) exists and \( r \leq \frac{3}{2} \). Detailed calculations show, that the finiteness of (3.11) and of \( \det U \) is probably a contradiction. So one would like to suppress \( \det U \). This is easily possible, because \( \det U \) plays the role of a renormalisation constant being for all integrals the same. It can be eliminated by defining the renormalized norm
\[ \| T_a(q) \|_{\text{ren.}} := \frac{\| T_a(q) \|}{\| T_0(q) \|} \]
where \( T_0(q) \) is the ground state functional. Then the renormalized norm is a completely finite number, provided the above mentioned conditions are satisfied. Still there remains a great arbitrariness concerning the choice of the weighting factors. This we try to remove in the next section.

4. Functional Scalar Products

For developing a quantum theory in functional space, we do not only need a norm definition for the state functionals themselves but also a functional scalar product between different state functionals. Such a scalar product can be given by a straightforward extension of the norm definition (3.4) resp. (3.14). Defining the weighted state functionals by
\[ W_a(q) := \exp \left\{ -\frac{1}{2} q G q \right\} \Gamma^{-\nu}(N) T_a(q) \]
where \( q \in G \) is a symbolic notation of the exponential in (3.4), the scalar product between two different state functionals \( T_a(q) \) and \( T_b(q) \) can be written
\[ \langle T_a(q) T_b(q) \rangle_{\text{ren.}} := \| T_0(q) \|^{-2} \int W_a(q) W_b(q) \text{d}r(q). \]
The existence of (4.2) for arbitrary state functionals \( T_a(q) \) and \( T_b(q) \) can be proven by the same method as used in section 3 and it is not necessary to repeat these statements. Then the norm definition is only a special case of the scalar product (4.2) and all these products are finite for physical state functionals. Therefore it follows that the weighted state functionals \( W_a(q) \) in (4.1) this equation can be formally written
\[ W_a(q) = S \cdot a \]
where \( S \) causes a mapping of the physical Hilbert space on to the functional Hilbert space. Contrary to the scalar product of different physical states \( \langle a | b \rangle \) in general the functional scalar product between corresponding functional states \( W_a(q) \) and \( W_b(q) \) is not orthogonal. This is due to the use of a weighting factor. Therefore \( S \) is in general a nonunitary similarity transformation. In this case it is convenient to introduce a second system, the so-called dual system of state functionals \( T^a(q) \) satisfying the orthonormality relations
\[ \langle T_a^a(q) T_b(q) \rangle_{\text{ren.}} = \delta_b^a. \]
Assuming the \( T^a(q) \) to be a linear combination of the original state functionals, we have
\[ T^a(q) = \sum_b c_{a b} T_b(q) \]
and by (4.4) follows immediately
\[ c_{a b} = \langle (T_a(q) T_b(q))_{\text{ren.}} \rangle^{-1}. \]
According to Schmidts orthogonalization procedure (4.6) exists if the vectors \( T_a(q) \) are linearly independent and normalizable. The linear independence is guaranteed by the linear independence of the \( | a \rangle \) states. The normalization can be performed...
by a proposal of Weber. Assuming for $G$ the ansatz
\[ G_{km} = \delta_{km} g_m \delta_{ab} \]  
(4.7)
the weighting factor contains just as many unknown parameters as there are normalization conditions for the different state functionals $T_a$. Assuming the norm expression for each functional to be unity by (3.8) one obtains a set of equations for the determination of the $g_m$. So in the framework of the norm and scalar product definition (4.2) we obtain a well defined dual set $T^\alpha$ connected with a non-unitary mapping of physical Hilbert space on to functional Hilbert space.

But now probably the definition (4.2) is not the only one that can be given for the construction of a functional Hilbert space and a mapping on to physical Hilbert space. Rather one has to expect an infinite set of possible mappings and functional Hilbert spaces. This gives rise to the question if there is any privileged functional Hilbert space or not. The answer is given in section 6. There it is shown that any functional Hilbert space leads to the same quantum mechanical information provided (4.4) is satisfied, i.e. the similarity transformation has to be a bijective one. So from the general physical point of view all possible functional Hilbert spaces are equivalent. This is true as long as one operates within the framework of quantum mechanics. But it will be shown in another paper that the special definition of the functional scalar product becomes physically important if one treats relativistic quantum field theory. Especially the definition given here is adapted to the functional formulation of nonlinear spinor theory of elementary particles.

5. Calculational Method for State Functionals

Being able to define a weighted norm for the state functionals, one would like to use this for the approximate calculation of these functionals. A method working with weighted norms is the method of least squares. This method has been applied successfully already in quantum mechanics and for Bethe-Salpeter problems. Unfortunately no rigorous proof of this method is known to the author, but one learns from quantum mechanics that for an application of this method the operator should be Hermitian at least and should have a lower limit of its spectrum. Now the eigenvalue equation (2.6) is formally Hermitian and its spectrum contains with the eigenvalue $\omega$ the eigenvalue $-\omega$ too, i.e. the spectrum is invariant against inversion at $\omega = 0$. Therefore by squaring of the operator no information is lost and the spectrum has a lower bound. So the conditions being supposed for the successful application of this method are satisfied for this problem. Writing the eigenvalue equation (2.6) in the symbolic form
\[ \mathbf{O} \left( \psi, \frac{\nabla}{\nabla} \right) + \omega \mathbf{T}(\psi) = 0 \]  
(5.1)
the condition of variational stationarity reads in this method
\[ \left\| \mathbf{O} \left( \psi, \frac{\nabla}{\nabla} \right) + \omega \mathbf{T}(\psi) \right\|^2 = \text{min} . \]  
(5.2)
This is always a positive expression. Therefore sequences of trial functions should give eigenvalues converging towards the exact values if the number of variational parameters is enlarged. For trial functions we use the approximate functionals
\[ T_N(\psi) := \sum_{n=1}^{N} \sum \frac{i^n}{n!} T_n(l_1, \ldots , l_n, a_1, \ldots , a_n, N) q_{l_1 s_1} \cdots q_{l_n s_n} \]  
(5.3)
where the free variational parameters are given by the expansion coefficients $T_n(l_1, \ldots l_n, N)$. To evaluate the equations for the trial functions (5.3) one derives first the exact equations for the state functionals from (5.2) and truncates afterwards these equations. We give only a short review what has to be done. Introducing superindices (5.1) can be written
\[ \sum_{m, n} \sum \frac{i^n}{n!} T_m(s_1, \ldots , s_m) O_{mn}(s_1, \ldots , s_m, k_1, \ldots , k_n) q_{k_1} \cdots q_{k_n} = 0 \]  
(5.4)

This representation follows from (2.6) by direct evaluation. It is the “integral” representation of (2.6).

By means of (5.4) and (5.5) the minimum expression (5.2) can be written

$$\sum_{m,j} T_{j}^{m} (s_{1} \ldots s_{j}) M_{jm}(s_{1} \ldots s_{j}, s_{1} \ldots s_{m}) T_{m}(s_{1} \ldots s_{m}) = \min$$

(5.6)

with

$$M_{jm}(s_{1} \ldots s_{j}, s_{1} \ldots s_{m}) = \sum_{n,1} \sum_{k_{1}, \ldots, k_{n}} O_{jm}^{n}(s_{1} \ldots s_{j}, k_{1} \ldots k_{n}) \frac{(-i)^{l}}{(l!)^{r}} \langle D_{l}(k_{1} \ldots k_{n} G) D_{n}(k_{1} \ldots k_{n} G) \rangle \frac{i}{(n!)^{r}} O_{nm}(s_{1} \ldots s_{m}, k_{1} \ldots k_{n}).$$

(5.7)

Then we consider $T$ and $T^{*}$ to be independent quantities and vary independently both of them. This gives

$$\sum_{m} \sum_{s_{1} \ldots s_{m}} M_{jm}(s_{1} \ldots s_{j}, s_{1} \ldots s_{m}) T_{m}(s_{1} \ldots s_{m}) = 0$$

(5.8)

and the conjugate complex equation, which brings no new information. As the scalar product of weighted Dyson functionals has been derived in App. I the expression of $M_{jm}$ can be calculated explicitly without any difficulty. Thus we do not evaluate the technic of solution in detail. It runs on the same pattern like in II and III.

### 6. Functional Quantum Theory

In the preceding sections we discussed all those operations which seem to be necessary to constitute a functional quantum theory i.e. a theory in functional space which is an isomorphism of ordinary quantum theory. To decide whether we have really obtained a functional quantum theory or not we consider the amount of physical information that can be derived by the functional methods and compare it with the maximum amount of information following from ordinary quantum theory. According to quantum theory the information is given by a complete set of quantum numbers, i.e. diagonalized observables and by the probability of the system to be in a certain state characterized by these numbers. Now for the anharmonic oscillator the only quantum number is the energy. It can be calculated either by conventional Schrödinger theory, or by functional calculus with equations (1.10) provided we have a powerful calculation method what we shall assume for this discussion. Therefore concerning the quantum numbers the statements of ordinary quantum theory and functional theory are the same and we have to look only for an equivalent probability definition in functional space.

To do this we remind of the probability definition in ordinary physical Hilbert space. In this space the most general state is defined by

$$\Psi(0) = \sum_{a} c_{a} \Psi_{a}(0)$$

(6.1)

and the occupation probability of a certain state $| a \rangle \equiv \Psi_{a}(0)$ is given by

$$P_{a}(0) = | c_{a} |^{2} = \langle \Psi(0) \Psi_{a}(0) \rangle.$$  

(6.2)

To imitate this in functional space we consider the generalized functionals

$$\Xi_{a}(j) := \langle 0 | T \exp i \int \psi_{a}(\xi) j_{a}(\xi) d\xi | \Psi(0) \rangle.$$  

(6.3)

Substitution of (6.1) into (6.3) then gives

$$\Xi_{a}(j) = \sum_{a} c_{a} \Xi_{a}(j)$$

(6.4)

where $\Xi_{a}(j)$ are the state functionals defined in (1.7), which is the analogue to (6.1). Transition to the rigged Hilbert representations changes (6.4) into

$$T(\xi) = \sum_{a} c_{a} T_{a}(\xi).$$

(6.5)

Assuming now the definition of a scalar product under the general conditions of section 4 and the existence of a corresponding dual set $T^{a}$ from (6.5) follows by using (4.4)

$$P_{a}(0) = | c_{a} |^{2} = \langle T_{a}(\xi) T^{a}(\xi) \rangle_{\text{ren}}.$$  

(6.6)

By this we see, that the probability statements of ordinary quantum theory can be reproduced in functional space. Therefore all physical information to
be obtainable from ordinary quantum theory can be reproduced by functional theory. Unfortunately the functional space is larger than the physical Hilbert space, and at present no selection principle between the desired and the undesired functionals is known. Nevertheless for the quantum theory of infinite degrees of freedom one may hope, that it is not necessary in future to work with the physical Hilbert space at all and quantum theory can be formulated independently in functional space only.

Acknowledgements

I thank very much Prof. Dr. W. Heisenberg and Dr. H. P. Dürr of the Max-Planck-Institut für Physik und Astrophysik, Munich, for their kind and critical reading of the manuscript. I am also very indebted to Dr. Schuler and Dr. Sohr, Dipl.-Phys. Dammaier, Dipl.-Phys. Scheerer and Dipl.-Phys. Weber for many discussion about the subject.

Appendix I

Already in III, App. I, II we discussed functional integration by means of the Hermitean functionals \(^2\)\(^1\). But the discussion given there is not very instructive because by the mathematical methods used in III the real structure of the problem is not revealed. A better structural understanding can be achieved by observing the isomorphism of the states for free fields with the Hermitean functionals. Although both quantities have a different meaning in connection with the present problem their formal definition is the same. As we perform our calculation in a \(L_2\)-representation we introduce from the beginning the \(L_2\)-variables, namely the \(Q_k\). By means of the \(Q_k\) the following creation and destructions operators can be defined

\[
A_k = 2^{-\frac{1}{2}}(Q_k + \frac{3}{2}\frac{1}{\partial Q_k}), \hspace{1cm} \bar{A}_k = 2^{-\frac{1}{2}}(Q_k - \frac{3}{2}\frac{1}{\partial Q_k}). \tag{1.1}
\]

Obeying the commutation relations

\[
[A_k \bar{A}_{k'}] = \delta_{kk'} 1 \tag{1.2}
\]

all other commutators vanish. Then the standard Hermitean functionals are defined by

\[
J_n(k_1 \ldots k_n) := \frac{1}{\sqrt{n!}} A_{k_1} \ldots A_{k_n} \exp -\frac{1}{2} \sum Q_k^2 \tag{1.3}
\]

and due to

\[
A_k \exp -\frac{1}{2} \sum Q_k^2 \equiv A_k |0\> = 0 \tag{1.4}
\]

the functional integral between two standard Hermitean functionals

\[
\langle J_l(k_1 \ldots k_l) J_n(k'_1 \ldots k'_n) \rangle = \int J_l(k_1 \ldots k_l) J_n(k'_1 \ldots k'_n) \prod dQ_k \tag{1.5}
\]

can be evaluated by the usual field theoretic methods giving

\[
\langle J_l(k_1 \ldots k_l) J_n(k'_1 \ldots k'_n) \rangle = \delta_{l,n} \frac{1}{n!} \sum_{i_1, \ldots, i_n} \delta(k_i - k'_i) \ldots \delta(k_n - k'_n). \tag{1.6}
\]

As long as the commutation relations (1.2) are satisfied these formulas are valid for usual indices as well as for superindices. Now from the definition of the functional variables \(Q_{kn}\) appearing in (2.1) it follows, that (1.2) is satisfied for superindices too. Therefore in the following all formulas can be read in indices as well as in superindices.

Further the connection between Dyson functionals and Hermite functionals is of interest. To derive it we introduce auxiliary variables \(x_k\). Then by (1.4) we have

\[
\exp\{\Sigma x_k A_k^2\} |0\> = \exp\{\Sigma x_k A_k^2\} \exp\{\Sigma x_k A_k\} |0\> \tag{1.7}
\]

and by application of the Hausdorff-formula on the right side of (1.7) follows

\[
\exp\{2\Sigma x_k Q_k\} |0\> = \exp\{\Sigma x_k A_k^2\} |0\> \exp\{\Sigma x_k \delta_{kk'} x_{k'}\}. \tag{1.8}
\]

Defining the standard Dyson functionals by

\[
D_n(k_1 \ldots k_n) := \frac{1}{n!} Q_{k_1} \ldots Q_{k_n} |0\> \tag{1.9}
\]

power series expansion in \(x_k\) and comparison of equal coefficients on both sides of (1.8) gives by observing (1.3) and (1.9)

\[
D_n(k_1 \ldots k_n) = \frac{1}{n!} \sum_{i_1, \ldots, i_n} \frac{1}{\sqrt{n!}} \frac{1}{(n-\mu)!} \frac{1}{\mu!} \frac{1}{2} \mu! \delta_{k_{i_1} k_{i_2}} \ldots \delta_{k_{i_{n-1}} k_{i_n}}. \tag{1.10}
\]
can be used for evaluating the scalar product of two Dyson functionals. This is not done here, because we are interested in more general definitions of Hermitean- and Dyson-functionals. They can be defined by introducing

\[ A_l(a) = U_l a U_l^{-1} \]

and their scalar product is given by

\[ \left \langle J_n(k_1 \ldots k_a) J_n(k_1 \ldots k_a) \right \rangle = \frac{1}{n!} \prod_{k=1}^{n} a_{k_1 k_2} \ldots a_{k_1 k_n} \det U . \]

Observing (1.6) and

\[ a_{kl} = \sum_r U_{kr} U_{rl} \]

(1.19) finally is given by

\[ \left \langle J_n(k_1 \ldots k_a) J_n(k_1 \ldots k_a) \right \rangle = \frac{1}{n!} \prod_{k=1}^{n} a_{k_1 k_2} \ldots a_{k_1 k_n} \det U . \]

In order to derive the connection between “weighted” Dyson functionals and weighted Hermite functionals, we introduce again auxiliary variables \( x_k \). Observing

\[ A^+_l(a) | 0 \rangle_w = 0 \]

and

\[ [A_l(a), A^+_k(a)] = a_{jk} \mathbf{1} , \]

we obtain in analogy to (1.7) and (1.8) the relation

\[ \exp \{ 2 \sum_k x_k a_{kl} | 0 \rangle_w \} = \exp \{ \sum_k x_k A^+_k(a) \} | 0 \rangle_w \cdot \exp \{ \frac{1}{2} \sum_{kl} x_k a_{kl} x_l \} . \]

Defining the weighted Dyson functional by

\[ D_n(k_1 \ldots k_a) := \frac{1}{n!} \sum_{l_1 \ldots l_a} \sum_{\mu} a_{l_1 k_1} \ldots a_{l_a k_a} J_n(k_1 \ldots k_a) a_{l_1 l_a}^{-1} \ldots a_{l_a l_a}^{-1} \]

power series expansion in \( x_k \) and comparison of equal coefficients on both sides of (1.24) gives by substitution of (1.13) and (1.25)

\[ D_n(l_1 \ldots l_a) = \frac{1}{n!} \sum_{i_1 \ldots i_n} \sum_{\mu} \frac{1}{(n-\mu)!} a_{l_1 k_1} \ldots a_{l_a k_a} J_n(k_1 \ldots k_a) a_{i_1 i_2}^{-1} \ldots a_{i_a i_a}^{-1} \]

By means of (1.26) and (1.21) the scalar product between two weighted Dyson functionals follows immediately. We obtain
\[ \langle D_n (l_1 \ldots l_n) \rangle D_m (l'_1 \ldots l'_m) \]
\[ = \frac{\det U}{n! m!} \sum_{\mu_1 \ldots \mu_n} C_{\mu_1 \ldots \mu_n}^{m} a_{l_{1 \mu_1}} \ldots a_{l_{n \mu_n}} a_{l'_{1 \mu_1}} \ldots a_{l'_{m \mu_m}} \]  
\[ \text{with } C_{\mu_1 \ldots \mu_n}^{m} := \left( \frac{1}{2} \right)^{m+n-\mu} \frac{1}{\mu! \left( \frac{m-\mu}{2} \right) ! \left( \frac{n-\mu}{2} \right)!}. \]

The scalar product between two standard Dyson functionals obviously is a special case of (I.27).

\section*{Appendix II}

For an estimate of (2.3) we make use of the definition
\[ T_n (a_1 \ldots a_n) := \int T_n (t_1 \ldots t_n) f_{u_1} (t_1) \ldots f_{u_n} (t_n) \, dt_1 \ldots dt_n \]  
and of the spectral representation of the $r$-functions by intermediate states
\[ r_n (x_1 \ldots x_n) = \sum_{m_1 \ldots m_{n-1}} \langle 0 | \psi_{a_1} | m_1 \rangle \ldots \langle m_{n-1} | \psi_{a_n} | a \rangle \exp \left\{ -i \sum_{j=1}^{n} (\omega_{m_j} - \omega_{m_{j-1}}) t_j \right\} \Theta (t_n - t_{n-1}) \ldots \Theta (t_1 - t_2) \]  
with $m_n = a$ in the exponential. By substitution of (II.2) into (II.1) follows
\[ T_n (a_1 \ldots a_n) = \int \exp \left\{ -i \sum_{j=1}^{n} x_j t_j \right\} \Theta (t_1 - t_2) \ldots \Theta (t_{n-1} - t_n) \, dt_1 \ldots dt_n \]  
with
\[ U(x_1 \ldots x_n, u_1 \ldots u_n) := \int \exp \left\{ -i \sum_{j=1}^{n} x_j t_j \right\} \Theta (t_1 - t_2) \ldots \Theta (t_{n-1} - t_n) \, dt_1 \ldots dt_n. \]  

Now from (II.4) follows the inequality
\[ |U(x_1 \ldots x_n, u_1 \ldots u_n)| \leq \prod_{j=1}^{n} \int |f_{u_j} (t_j)| \, dt_j \]  
and for the integrals an estimate can be given
\[ \int |f_{u_j} (t_j)| \, dt \leq C \]  
with a finite constant $C$ for arbitrary $u$ by means of an inequality for Hermitian polynomials. Therefore (II.4) is uniformly bounded over the whole range of the variables $x_1 \ldots x_n, u_1 \ldots u_n$ and satisfies the inequality
\[ |U(x_1 \ldots x_n, u_1 \ldots u_n)| \leq C \]  
for all $n$. For an estimate of the first part in (II.3) we first observe, that according to the properties of matrixfactors for the oscillator functions (1.5) the equation

\[ 25 \quad \text{I. M. Ryshik and I. S. Gradshteyn, Tafeln, Deutscher Verb. d. Wiss., Berlin 1957, ge. 7.124.} \]