Quantum field propagator for extended-objects in the microcanonical ensemble and the S-matrix formulation.

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Abstract
Starting with the well-known Nambu-Goto action for an N-extended body system the propagator in the microcanonical ensemble is explicitly computed. This propagator is independent of the temperature and, in contrast with the previous references, takes account on all the non-local effects produced by the extended objects (e.g., strings) in interaction. The relation between relativistic quantum field theories in the microcanonical approach and the pure S-matrix formulation is established and analyzed.

Keywords: Microcanonical ensemble, S-matrix formulation, string propagators.
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1 Introduction

Several studies into the behaviour of relativistic quantum field theories at finite temperature have advanced considerably since the initial papers in 1974 [1-3]. In addition to the imaginary time formulation of Matsubara, two
more methods are available for formulating real time perturbation theory with temperature dependent propagators: the functional integral approach [5], and the operator approach of [6]. The calculations have all been realized in the canonical ensemble, specified by volume $V$ held in contact with a heat bath at a fixed temperature $T$. The partition function and the scalar propagator were defined by [7]

$$Z_{\beta} = \sum_n \exp(-\beta E_n)$$

$$Z_{\beta} D_{\beta}(t, x) = -i \sum_n \exp(-\beta E_n) \langle n | T \left[ \phi(t, x) \phi(0, 0) \right] | n \rangle$$

where $|n\rangle$ label a complete set of energy eigenstates and $\beta = 1/T$, as usual. Since these propagators have singularities whose locations depend on the temperature, some inconsistencies appear: when the temperature is non zero the energy differences that are independent of $T$ lead to singularities whose locations are, apparently, $T$-dependent. To solve this puzzle, in [8] it was shown that the natural context to study the same field theory is in the microcanonical ensemble where the system is completely isolated at volume $V$, then the total energy of the system $E$ will remain constant. The microcanonical partition function and propagator were naturally defined as

$$\overline{Z}_{E} = \sum_n \delta(E_n - E)$$

$$\int_0^E dE' \overline{Z}_{E-E'} \overline{D}_{E'}(t, x) \equiv -i \sum_n \delta(E_n - E) \langle n | T \left[ \phi(t, x) \phi(0, 0) \right] | n \rangle$$

The mapping between the ensembles is realized by Laplace transform from $E$ to $\beta$. This distribution is more physically compelling in applying quantum field theory to the early universe or to heavy-ion collisions since the systems are completely isolated and are not certainly in contact with a heat reservoir.

By the other hand, in ref.[19] the connection between dual amplitudes (i.e., interacting strings) and thermodynamics properties of a strongly interacting system was studied in the context of the canonical approach of ref.[18]. As a result from this study, the authors of [19] showed that the strong duality is loss in this strongly interacting system: the narrow resonances aproximation does not reproduces the Regge model results.
In this work, strongly motivated for above arguments, we give the next step in studying relativistic quantum field theories in the context of the microcanonical formulation: firstly, establish the relation with the S-matrix formulation following a similar procedure as in refs. [12,18] in the context of the canonical ensemble; and secondly, compute the microcanonical propagator for a geometrical non-local and nonlinear action that is the well-known Nambu-Goto action for a system of strings. The plan of this paper is as follows: in Section 2, we describe the microcanonical formulation briefly. In Section 3, the relation between the microcanonical approach in relativistic quantum field theories and the axiomatic S-matrix formulation is established. Finally, in Section 4, the microcanonical propagator for an N-extended body system described by the Nambu-Goto action is successfully performed and analyzed.

2 Microcanonical formulation

The true vacuum for a statistical system in thermal equilibrium can be obtained writing the termal vacuum in terms of the density matrix \( \hat{\rho} \)

\[
|0(\beta)\rangle = \hat{\rho}(\beta, H)|\mathcal{J}\rangle,
\]

where

\[
\hat{\rho}(\beta, H) = \frac{\rho(\beta, H)}{\langle \mathcal{J} | \rho(\beta, H) | \mathcal{J} \rangle}
\]

\[
\rho(\beta, H) = e^{-\beta H}
\]

\[
|\mathcal{J}\rangle = \left[ \prod_{k,m} \sum_{n_{k,m}} \prod_{k,m} |n_{k,m}\rangle \otimes |\tilde{n}_{k,m}\rangle \right]
\]

The trace of an observable operator is given by

\[
Tr\hat{O} = \langle \mathcal{J} | \hat{O} | \mathcal{J} \rangle
\]

For example, the free field propagator can be determined from

\[
\Delta^{ab}_{\beta} = -i \langle \mathcal{J} | T \phi^a(x_1) \phi^b(x_2) \hat{\rho} | \mathcal{J} \rangle,
\]
where the superscripts on $\phi$ refer to the member of the thermal doublet being considered (for details see ref. [9])

$$\phi^a = \begin{pmatrix} \phi \\ \tilde{\phi}^\dagger \phi \end{pmatrix}$$

As was shown in [9], the Fourier transform of $\Delta_{11}^\beta (x_1, x_2)$, that is the physical component, is equal to

$$\Delta_{11}^\beta = \frac{1}{k^2 - m^2 + i\epsilon} - 2\pi i \delta (k^2 - m^2) n_{11}^\beta (m, k)$$

(11)

where in a canonical ensemble; for a thermal system at temperature $\beta^{-1}$ the number density is given by $n_{11}^\beta (\omega) = \frac{1}{e^{\omega/\beta} - 1}$.

If instead of treating the $n$-body system as objects in thermal equilibrium at fixed temperature $T$ and corresponding vacuum $|\beta H\rangle$ we treat the system as having fixed energy $E$, we can formally define the microcanonical vacuum as [10]:

$$|E\rangle = \frac{1}{\Omega (E)} \int_0^E \Omega (E - E') L_{E-E'}^{-1} \Omega (E) |\beta H\rangle dE'$$

(12)

where $L^{-1}$ is the inverse Laplace transform. Using this basis, the physical correlation functions are expressed as

$$G_{\alpha_1,\ldots,\alpha_N}^a (1, 2, \ldots, N) = \langle J| T \phi^{\alpha_1} (1), \ldots, \phi^{\alpha_N} (N) |E\rangle$$

(13)

The physical component of the propagator in the microcanonical field formulation [9,10] is given by

$$\Delta_{11}^E (k) = \frac{1}{k^2 - m^2 + i\epsilon} - 2\pi i \delta (k^2 - m^2) n_{E} (m, k) ,$$

(14)

where the second term in the above expression corresponds to the statistical (microcanonical) part and

$$n_{E} (m, k) = \sum_{l=1}^{\infty} \frac{\Omega (E - l\omega_k (m))}{\Omega (E)} \theta (E - l\omega_k)$$

(15)

is the microcanonical number density, where $l, \omega_k$ and $\theta (E - l\omega_k)$ are the mode number, the dispersion relation and the step function, respectively.
3 Axiomatic S-matrix formulation in QFT: microcanonical description

The relation between the S matrix formulation in QFT and the statistical operator in the canonical ensemble [11,12,18] can be easily extended to the microcanonical description as follows. The propagator operator (we assume it with its statistical part)

\[ \hat{G}(E) = \frac{1}{E - \hat{H} - i\epsilon} + G_{st}, \]  

(16)
namely, its imaginary part (where \( G_{st} \) is the statistical part of the full propagator) and the S-matrix are related as

\[ \Im(\hat{G}(E)) = \Im(\hat{G}_0(E)) + \frac{1}{4i} \hat{S}^{-1}(E) \frac{\partial}{\partial E} \hat{S}(E) \]  

(17)

where \( \hat{S}(E) \) is the scattering operator at the energy \( E \) and \( \hat{G}_0(E) \) is the free part of the full propagator (that leads to non-connected graphs in the cluster expansion). The relation to the physical \( \hat{T}(E) \) matrix is

\[ \hat{S}(E) = 1 + i\delta(E - \hat{H}_0) \hat{T}(E), \]  

(18)

\( \hat{H}_0 \) being the free Hamiltonian. The logarithm of the trace of \( Z(T,V) \) can be written as [11]

\[ \ln Z(T,V) = \ln Z_0(T,V) + Tr \int e^{-\bar{b}_\mu P_\mu} \Delta^{3} \left( \bar{P} - \hat{P} \right) \Im \left[ \hat{G}(E) - \hat{G}_0(E) \right] \right]_{\text{connected}}, \]  

(19)

where the four-vector temperature \( \bar{b}^\mu \) is defined by the identity

\[ \bar{b}^\mu = \frac{1}{T} u^\mu \]  

(20)

\( T \) is the temperature in the rest frame of the box enclosing the thermodynamical system, \( u^\mu u_\mu = 1 \) being its four-velocity. The index \( \text{connected} \) means that only the connected part of the full propagator \( \hat{G}(E) \) must be taken: simply by substracting to the full propagator the free part \( \hat{G}_0(E) \) corresponding to an ideal gas configuration.
Recalling that the operation of taking the connected part leaves the operations invariant, we can write

$$\ln Z (T, V) = \ln Z_0 (T, V) + \int e^{-\beta P} d^4 P \rho_I (P^2, V.P, V^2)$$  \hspace{1cm} (21)

$$\forall \mu \equiv \frac{2V}{(2\pi)^3} u^\mu \text{ (four-vector volume)},$$  \hspace{1cm} (22)

where, by analogy with the interaction level density, the function which we call mass spectrum, or cluster level density, is given by

$$\rho_I (P^2, V.P, V^2) = \left\{ -\frac{1}{\pi} \delta (P - \hat{P}_0) \Im \left[ \hat{G} (E) - \hat{G}_0 (E) \right] \right\}_{\text{connected}}$$  \hspace{1cm} (23)

here the index $I$ means "interaction". On the other hand, from the pure statistical point of view,

$$\rho_I (P^2, V.P, V^2) = \rho - \rho_0,$$  \hspace{1cm} (24)

$\rho$ being the full density of states of the system under consideration and $\rho_0$ the density of states in the free configuration, explicitly

$$\rho_0 = \sum_{k=1}^{\infty} \forall \mu P_\mu \delta \left( P^2 - k^2 m^2 \right)$$  \hspace{1cm} (25)

$\rho$ satisfies the following relation:

$$\rho (P^2, V.P, V^2) = \forall \mu P_\mu \rho (\sqrt{P^2})$$  \hspace{1cm} (26)

The above condition is important. This means that the cluster counting $\rho (P^2, V.P, V^2)$ in the rest frame of the system can be reexpressed by the counting of states for a single particle of the $\rho (m)$ mass degeneracy moving in the volume $V$. Now, since the mathematical mapping between the microcanonical and canonical ensembles is easy to see, the $\Omega (E, V)$ we looking for is

$$\ln \Omega (E, V) = \ln \Omega_0 (E, V) + Tr \int d^4 P \left\{ -\frac{1}{\pi} \delta^3 (P - \hat{P}) \Im \left[ \hat{G} (E) - \hat{G}_0 (E) \right] \right\}_{\text{conn}}$$  \hspace{1cm} (27)
Notice that in order to obtain the above result, the main difference of our procedure with the procedure given in refs.[12,18] for the canonical ensemble is that we take as a starting point the propagator operator eq.(16), with its statistical part. Considering that the imaginary part $\Im \left[ \hat{G}(E) - \hat{G}_0(E) \right]$ is the connected part of the full statistical propagator $\hat{G}(E)$, it corresponds in the microcanonical field formulation to the physical component $\Delta_{E}^{11}(k)$ eq.(14), which leads to an analog expression like (27), we can easily seen from the comparison between the propagator, eq.(16), with expressions, eq.(23-26), yielding the level density of states $\rho_I$ that the interactions between particles in the system (dynamical information from the connected part of the propagator (16)) are automatically translated into variation of the energy (mass) levels in the statistical ensemble (given by $\rho(m)$ and the statistical information of the propagator (16)). That is, the study of any simple interactions between particles is equivalent to that of the thermodynamical properties of the statistical ensemble of such particles as a whole.

4 The Nambu-Goto action and the microcanonical propagator

It is difficult to study this system in the Hamiltonian framework because of the constraints and the vanishing of the Hamiltonian. As is known, the Nambu-Goto action is invariant under the reparametrizations

$$\tau \to \tilde{\tau} = f_1(\tau, \sigma) \quad \sigma \to \tilde{\sigma} = f_2(\tau, \sigma)$$

then, we can make the following choice for the dynamic variable $x_0$ and the space variable $x_1$, as first proposed by B. Barbashov and N. Chernikov in ref.[13] which do not restricts the essential physics and simplifies considerably the dynamics of the system

$$x_0(\tau, \sigma) \equiv x_0(\tau) \quad x_1(\tau, \sigma) \equiv \kappa \sigma \quad (\kappa = \text{const})$$

for this, it is sufficient to use the chain rule of derivatives and to write the action in the form

$$S = -\frac{\kappa}{\alpha} \int_{\tau_1}^{\tau_2} \dot{x}_0 d\sigma \, d\tau \sqrt{\left[1 - (\partial_0 x_b)^2\right] \left[1 + (\partial_1 x_a)^2\right]}$$
\( a, b = 2, 3; \partial_1 x_a = \varepsilon_{1a}^0 \partial_0 x_b, \) where in order to simplify at maximum this action we choose an orthonormal frame. (Thus, we pass from the Nambu-Goto action to the Born-Infeld representation).

Therefore, the invariance with respect to the invariance of the coordinate evolution parameter means that one of the dynamic variables of the theory \((x_0(\tau)\) in this case) becomes the observed time with the corresponding non-zero Hamiltonian

\[
H_{BI} = \Pi_a x^a - L = \sqrt{\alpha^2 - \Pi_b \Pi^b}, \tag{28}
\]

where

\[
\Pi^b = \frac{\partial L}{\partial (\partial_0 x_b)}, \quad \alpha \equiv \frac{\kappa}{\alpha'} \sqrt{1 + (\partial_1 x_a)^2}
\]

Now in order to find the free propagator from the NG Hamiltonian we proceed as follows:

From the simplest quantum path-integral formalism, we have

\[
K(q', t, q, 0) \equiv \langle q' | (e^{H\varepsilon})^N | q \rangle = \langle q' | \Psi(r, s, t...) \rangle,
\]

where \(K(q', t, q, 0)\) is the propagator, \(H\) is the Hamiltonian of the theory \(t\) is the time that was fractionated in small lapses \(t = N\varepsilon\) and \(q, q'\), and \(\Psi(r, s, ...)\) are the physical states with \(r, s...\) quantum numbers. With the usual path integral operations and introducing the integral representation for a pseudodifferential operator [14]

\[
\int (t^2 + u^2)^{-\lambda} e^{itx} dt = \frac{2\pi^{1/2}}{\Gamma(\lambda)} \left( \frac{|x|}{2u} \right)^{\lambda-1/2} K_{\lambda-1/2} (u |x|)
\]

where \(K_\nu(x)\) is the MacDonald’s function, the propagator for a sub-interval takes the form

\[
K_{q_j, q_{j+1}} = \delta_{q_j, q_{j+1}} - i\varepsilon \left[ 4\alpha K_{-1} (\alpha |q_j - q_{j+1}|) \right] \left( \frac{q_j - q_{j+1}}{|q_j - q_{j+1}|} \right)
\]

Putting on all the subinterval propagators together, we obtain the full propagator

\[
K = \delta_{q_N, q_0} - iN\varepsilon \left[ 4\alpha K_{-1} (\alpha |q_N - q_0|) \right] \left( \frac{q_N - q_0}{|q_N - q_0|} \right)
\]
Making, without loss of generality, the transformation $-it \rightarrow -\beta$, $q_0 = q_N$ integrating and Fourier transforming to momentum space, yields the canonical partition function $Z_c$

$$Z_c = \sum_N \left\{ 1 - \beta \left[ 4\alpha \frac{K_{-1} (\alpha |q_N - q_0|)}{|q_N - q_0|} \right] \right\}$$

$$= \sum_N \exp \left\{ -\beta \left[ 4\alpha \frac{K_{-1} (\alpha |q_N - q_0|)}{|q_N - q_0|} \right] \right\}$$

The microcanonical partition function $\Omega_m$ is obtained as the inverse Laplace transform of the last expression:

$$\Omega_m = \delta (E) - \sum_{N=1}^{\infty} \sum_{n_1=1}^{\infty} \cdots \sum_{n_N=1}^{\infty} \left[ 4\alpha \frac{K_{-1} (\alpha \sum n_j \varepsilon_j - E)}{E^2} \right] \frac{1}{n_1 n_2 \cdots n_N}$$

where the factor $\frac{1}{n_1 n_2 \cdots n_N}$ allows one to eliminate the overcounting, and $\sum n_j \varepsilon_j = E_N$.

The free microcanonical propagator for the N-extended body system can be consistently formulated using the relation (valid for free fields) between time ordered products and normal products

$$-iT [\varphi (x) \varphi (0)] = D_F (x) - i : \varphi (x) \varphi (0) :$$

where $T$ is the temperature of the thermal bath, $D_F (x) = -i \langle J | \varphi (x) \varphi (0) | J \rangle$ is the ordinary Feynman propagator with the expectation value evaluated in the basic states of our system (i.e., for zero temperature)

$$|J\rangle = \left[ \prod_{k,m} \sum_{n_k,m} \prod_{k,m} [n_{k,m}] \otimes |\tilde{n}_{k,m}\rangle \right]$$

Since the relation between microcanonical-canonical formulations is via a Laplace transform, it is reasonable to perform the following mapping:

$$\int_0^\infty dE' \Omega_E - E' D_{E'} =$$

$$= D_F \Omega_E - i \sum_{N=1}^{\infty} \sum_{n_1=1}^{\infty} \cdots \sum_{n_N=1}^{\infty} \left[ 4\alpha \frac{K_{-1} (\alpha \sum n_j \varepsilon_j - E)}{E^2} \right] \frac{1}{n_1 n_2 \cdots n_N} \langle J | : \varphi (x) \varphi (0) : |E \rangle,$$

(29)
where we defined the microcanonical state as

$$|E⟩ = \frac{1}{Ω_E} \int_0^∞ dE′Ω_{E-E′}L_{E′}^{-1}[|β⟩],$$

$L^{-1}$ being the inverse Laplace transform and $D_{E′}$ being the microcanonical propagator. The matrix element for the most general states in our system is

$$⟨J| : \varphi(x) \varphi(0) : |E⟩ = \frac{1}{Ω_E} Ω_{E-E′} \sum_{p} \left[ \frac{n}{ε_pV} \cos (p.x - ε_p t) \right]$$

By inserting this expression in the definition of the microcanonical propagator $D_{E′}$ (39) given above and converting the momentum sum into an integral, we have

$$D_{E}(t, x) = δ(E) D_{E} - 4iα \int \frac{d^3p}{(2π)^3} \sum_{n=1}^∞ \frac{K_{-1}(α|n_jε_j - E|)}{E^2} \frac{Ω(E-n_jε_j)}{Ω(E)} \cos (p.x - ε_j t)$$

Finally, Fourier transform to momentum representation gives

$$D_{E}(t, x) = \frac{δ(E)}{ω^2 - k^2 - m^2 + iε} - 8πiαδ(ω^2 - k^2 - m^2) \sum_{l=1}^∞ \frac{K_{-1}(α|lω_k - E|)}{E^2} \frac{Ω(E-lω_k)}{Ω(E)} θ(E-lω_k),$$

(30)

where $θ(x)$ is the usual step function. The first term in the microcanonical propagator is the usual (non-termal) Feynman propagator, the second one is the new microcanonical statistical part. This part is crucial for the correct description of the full N-extended body system, as we can see explicitly expanding the Mac Donald’s function $K_{-1}$ in the second term of the free microcanonical propagator

$$8πiαδ(ω^2 - k^2 - m^2) \sum_{l=1}^{M/ω_k} \left[ \ln \left( \frac{γ_ε|ω_k-E|}{2} \right) \frac{α|ω_k-E|}{2} \sum_{s=0}^∞ \frac{α|ω_k-E|/2^{s}}{sΓ(s+1)} - \right.$$ 

$$\left. \frac{1}{2} \sum_{s=0}^∞ \frac{(α|ω_k-E|/2)^{2s+1}}{sΓ(s)} \left( \sum_{h=1}^∞ \frac{1}{h} + \sum_{h=1}^∞ \frac{s+1}{h} \right) + \frac{α|ω_k-E|}{4} \right] \frac{θ(E-ω_k)}{Ω(E)} \frac{Ω(E-ω_k)}{Ω(E)}$$

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and expressing it as a function of the mass, \( M \) being the energy \( E \) of the extended-bodies (strings) in our system \( (E = M \equiv \text{total mass of the system}) \)

\[
8\pi i\alpha\delta (\omega^2 - k^2 - m^2) \sum_{l=1}^{M/\omega_k} \left[ \ln \left( \frac{2\omega_k - M}{2} \right) \right] - \frac{1}{2} \sum_{s=0}^{\infty} \frac{(\alpha|\omega_k - M|/2)^{2s+1}}{s!\Gamma(s)} \left( \sum_{h=1}^{s} \frac{1}{h} + \sum_{h=1}^{s+1} \frac{1}{h} \right) + \frac{\alpha|\omega_k - M|}{4} \right] \frac{\theta(M - \omega_k) \Omega(M - \omega_k)}{M^2 \Omega(M)}
\]

We see that when \( M \to 0 \) this expression yields the pure string-like behaviour (Gamma type string-amplitude). Notice that previously \([15,16]\) the relation between the Feynman propagator and the Veneziano amplitude was put "by hand". We see that this type of structure coming from the statistical microcanonical part is contained in our microcanonical propagator, eq.(30). It should be noted that the relation between temporal and normal ordering of the field operators contributes to the statistical part of this full propagator. These implications will be discussed everywhere \([17]\), where we will focus on some concrete problems (arrow of time, the early universe, etc.).

It is interesting to note that the main difference between our microcanonical propagator, eq.(30), and the \( \Delta_{11} \) of the previous section is in that the propagator, eq.(30), includes all nonlocal effects: from the \( n \)-bodies of the system as extended objects (i.e., strings) and from the derivation of this propagator from a theory with a Hamiltonian not quadratic in momenta as in the Nambu-Goto formulation of the string theory. For instance, the propagator, eq.(30), becomes the propagator, eq.(14), when the extended bodies (e.g., strings) became point-particles and the Hamiltonian is quadratic in momenta (constrained particle Hamiltonian).

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