

Time-symmetric relativistic and quantum theories. 3. Time-symmetric QFT

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Annotation

Time-symmetric relativity theory (TSR) based on Zisman-Stueckelberg-Feynman (ZSF) interpretation includes reference frames going backward in time, maximally expanding their symmetry group to general Lorentz group $O(1,3)$ with 4-inversion (Article 1). Similar time-symmetric (TS) extensions of relativistic quantum mechanics (TS RQM) (Article 2) and QFT (TS QFT) (this article) avoid problems of former RQM and QFT. Former QFT introduced unfounded hypotheses about physical quantum fields in entire space for each type of particle, and these hypotheses led to insoluble problems. TS QFT is limited to known external interaction fields, and describes particles using theory of many-particle systems. In second quantization method, single-particle wave functions and operators, which increase or decrease number of particles in system, form operator amplitudes of transition probabilities. These amplitudes are complex, act in their state spaces for given energy sign, and lead to positive state norm. As in quantum mechanics, they, together with conjugate amplitude, form canonical pair. Canonical formalism in relativity theory is consistent at using Hamilton-Jacobi (HJ) formalism, where 3-momentum and Hamiltonian enter into HJ equation in the same degree. Differential form of relativistic HJ equation is Klein-Gordon (KG) equation, expressing the relationship between energy and momentum, and is general equation for physical components of all operator amplitudes. For bosons, KG equation suffices; for fermions, linearization of HJ equation is required, leading to special first-order equations, such as Dirac equation. Operators of observables are normally ordered, and there are no zero-point fluctuations in ground states. Probability current is determined by 4-momentum and has the same sign; interaction currents are also proportional to probability current and change sign with 4-momentum, which corrects the currents in fermion theory. Introducing interactions into HJ equation leads to new contact terms. Time-symmetric orderings of operator amplitudes lead to causal propagators and usual diagram technique. Taking into account external gravitational field of quanta at Planck length, gravitational radius of the system, and dilation of their proper times in this field lead to gravitational self-regularization of loop diagrams. This makes perturbation theory series in Standard Model (SM) and quantum gravity finite and convergent. Contributions of chiral anomalies are finite and cancel out in SM.

Key words: antiparticles, negative energy particles, Zisman-Stueckelberg-Feynman interpretation, general Lorentz group, 4-inversion, quantum fields, zero-point vacuum energy

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1. Introduction

In the covariant formulation of relativistic quantum theory, there are states of negative energy associated with the existence of antiparticles. This connection is nontrivial and has led to a number of interpretations over nearly a century.

The first attempt to give them a physical meaning was Dirac's famous hole theory [1]. In this hypothesis, the proper time τ of negative energy particles, in contradiction with relativity theory, was considered to go forward along with ordinary time t .

A consistent description of negative energy states in relativity theory was given by Zisman (1940) [2], Stueckelberg (1941) [3] and Feynman (1949) (ZSF) [4]. They showed that in relativity theory, negative energy particles travel backward along the axis t only, i.e., $\Delta\tau > 0$ at $\Delta t < 0$, and their description is equivalent to the description of positive energy antiparticles traveling forward in time. The interpretation of negative energy states as a way of describing the states of antiparticles, ZSF interpretation, underlies relativistic quantum mechanics (RQM) and its covariant diagram techniques.

However, relativity theory itself was not fully consistent with the interpretation of the ZSF, since it did not include the rest frames of negative-energy particles, where their proper time and rest energy should be measured [5,6]. In the first article [7], relativity theory was extended to such reference frames going backwards along the axis t , formulating the time-symmetric relativity theory (TSR) and extending the symmetry group to the general Lorentz group $O(3,1)$ including time inversion, and 4-inversion $x_{-}^{\mu} = -x^{\mu}$, $\mu = 0, 1, 2, 3$ in general case.

TSR led to a time-symmetric RQM (TS RQM) free of problems of former RQM (Article 2 [8]). The probability current in TS RQM is determined by 4-momentum, and their signs coincide; therefore, probabilities are positive, and the sign of the current depends on the direction of temporal evolution. Interaction currents, if proportional to the probability current, change sign when 4-momentum changes sign, which is the case for interactions with gauge fields. All this made the theory of bosons consistent and introduced corrections to the theory of fermions.

In this third article, the application of TSR to QFT will be considered. The former form of QFT postulated the existence of physical quantum fields filling all space for each type of particle [9-11]. However, in reality, these hypotheses were unfounded and only led to insoluble problems. One of the main ones was the problem of zero-point vacuum fluctuations. For the electromagnetic field, this consequence of the quantum field was refuted experimentally. The Casimir effect showed that only the effects of field fluctuations of real sources, the atoms of the crystal, are observed, and there are no

consequences of zero-point fluctuations of the electromagnetic field vacuum, which should double the effect. The same situation applies to the contribution of zero-point fluctuations to the Lamb shift and anomalous magnetic moments – the experiment shows the contribution of only loop corrections, the contribution of real sources, and the absence of the contribution of zero-point fluctuations of the field vacuum, which should also double the observed effects.

In the article the time-symmetric QFT (TS QFT) based on TSR is formulated, free of the hypotheses and problems of the former QFT. It retains the formalism of QFT, not because its hypotheses were correct, but because their undesirable consequences were manually corrected so that the formalism corresponds to experimental results. In TS QFT, the physical fields are considered to be the external interaction fields known from practice, and hypotheses about quantum fields are excluded. The quantum properties of external fields are reduced to the emission and absorption of their quanta, which, like other particles of matter, are further described by the methods of the theory of many-particle systems, which conveniently introduces effective fields.

In one of them, the second quantization method, field operators are introduced as products of single-particle wave functions and operators that increase or decrease the number of particles. Their physical meaning is that they are operator transition amplitudes and therefore complex, acting in their state spaces for each energy sign and yielding positive state norms. Like wave functions in quantum mechanics, these operator amplitudes form a canonical pair with their adjoint operator. Further, the description of antiparticles in terms of negative-energy particles traveling backward in time makes TS QFT covariant and consistent.

In the relativistic relation between energy and momentum, both of them enter quadratically, and in the canonical formalism, this property is consistently taken into account in the Hamilton-Jacobi (HJ) method [8]. The differential form of the relativistic HJ equation, the Klein-Gordon (KG) equation, expressing the relation between energy and momentum, is the general equation for physical components of the operator transition amplitudes of all particles. In the case of bosons, the general KG equation is sufficient, and for fermions, the linearization of the HJ equation leads to special first-order equations, in particular, to the Dirac equation. Hamiltonians in TS QFT do not contain the zero-point energy of the ground state and lead to normally ordered operators of observables. The introduction of interactions into the HJ equation leads to new contact terms, which lead to observable corrections.

The time-symmetric ordering of the operator transition amplitudes leads to causal propagators and to the covariant diagram technique, and thus TS QFT derives from physical principles the former heuristic recipes that allowed the processes of particle physics to be correctly described.

As in TS RQM, the external gravitational field generated by the total energy of the system of particles in the loop diagram is taken into account. Accounting for the slowing down in this external field of the particles' proper times relative to world time, in which the S-matrix is determined, leads to the freezing of all processes at the Planck length, the gravitational radius of the system of particles [12]. This gravitational self-regularization makes the perturbation series of the Standard Model (SM) and quantum gravity both finite and convergent. Chiral anomalies grow linearly and their contributions are large, but their cancellation in SM becomes a rational operation.

In Section 2 TS RQM and the transition from it to TS QFT are discussed. In Sections 3–4 time-symmetric theories of effective boson and fermion quantum fields, their commutators and propagators are considered. In Section 5 the theory of interacting particles and perturbation theory, as well as gravitational self-regularization of loop diagrams are shortly discussed. TS QFT is presented in more detail in the book [13]. In the article $c = \hbar = 1$ and the metric $g_{\mu\nu} = \text{diag}(1, -1, -1, -1)$.

2. From time-symmetric RQM to time-symmetric CFT

2.1. Canonical structure of the state space of quantum theory

In quantum mechanics, states are described by complex wave functions $\psi(\mathbf{x}, t)$, $\psi^*(\mathbf{x}, t)$, which are probability amplitudes and are related with the probability density as $\psi^* \psi = \rho$. They form a canonical pair in phase space (ψ, ψ^*) with Poisson brackets:

$$\{A, B\} = \frac{1}{i\hbar} \int \left(\frac{\delta A}{\delta \psi} \frac{\delta B}{\delta \psi^*} - \frac{\delta B}{\delta \psi} \frac{\delta A}{\delta \psi^*} \right) d^3\mathbf{x}, \quad (1)$$

$$\{\psi(x, t), \psi^*(x', t)\} = \delta(x - x') / i\hbar, \quad (2)$$

$$\{\psi(x, t), \psi(x', t)\} = \{\psi^*(x, t), \psi^*(x', t)\} = 0.$$

The algebra of observables in this phase space leads to an operator algebra. The values of the observables A are defined through bilinear forms of amplitudes with an operator $\hat{A}$ acting on these states:

$$A = \int \psi^* \hat{A} \psi d^3\mathbf{x}. \quad (3)$$

The corresponding evolution equation for the observable is:

$$\frac{\partial A}{\partial t} = \{A, H\}, \quad (4)$$

where H is the Hamiltonian, also defined by a bilinear form of the form (3). The canonical equations in this phase space for the canonical pair itself are:

$$\frac{\partial \psi}{\partial t} = \{\psi, H\} = \frac{1}{i\hbar} \frac{\delta H}{\delta \psi^*} = \frac{1}{i\hbar} \hat{H} \psi, \quad (5)$$

$$\frac{\partial \psi^*}{\partial t} = \{\psi^*, H\} = -\frac{1}{i\hbar} \frac{\delta H}{\delta \psi} = -\frac{1}{i\hbar} \hat{H} \psi^* \quad (6)$$

and represent the Schrödinger equations.

This canonical structure of the state space is preserved in any application of the quantum mechanical formalism, including TS RQM and TS QFT. In TS RQM, this is unquestionable, since the physical meaning of the wave function remains the same, thus below only the main results will be presented.

The situation in TS QFT is more nontrivial, since the states are multiparticle, and the coefficients of their momentum expansions in single-particle wave functions ψ, ψ^* become operators that lower or raise the number of particles in a given state. Probability amplitudes therefore become effective field operators. Probabilities are still defined through their bilinear forms, and the former wave equations naturally remain be equations for these operator transition amplitudes. Therefore, changing the mathematical form of the canonical pair to an operator form does not alter the former structure of the phase space. What changes is that it becomes the canonical structure in the phase space of operator transition amplitudes $(\hat{\psi}, \hat{\psi}^+)$, where Hermitian conjugation is introduced (henceforth, we will omit the hats).

2.2. Time-symmetric relativistic quantum mechanics

TS RQM (see Article 2 [8]) retained the fundamental results of the former RQM, expressed in wave equations and diagram techniques, while changes their derivations and interpretations, which led to many problems. As a result, the theory became more consistent, eliminating a number of former contradictions and problems. Below, we will consider those results of TS RQM that must be taken into account at constructing TS QFT.

The expressions for the Hamiltonian of a free relativistic particle in terms of 3-momentum are quadratic and therefore, relativistic HJ equation, as well as the relativistic Schrödinger equation, also contain these operators in quadratic form:

$$-\frac{\partial^2 \psi}{\partial t^2} = H_0^2 \psi, \quad H_0^2 = \mathbf{p}^2 + m^2. \quad (7)$$

Substituting into (7) the expression for 3-momentum $\mathbf{p}^2 = -\nabla^2$ gives the KG equation:

$$(\partial_\mu \partial^\mu + m^2) \psi = 0. \quad (8)$$

Linearization of the quadratic Hamiltonian in (7) gives the matrix Hamiltonian:

$$H_0 = \boldsymbol{\alpha} \mathbf{p} + \beta m. \quad (9)$$

Substituting (9) into the Schrödinger equation (5) turns it into the Dirac equation:

$$(i\gamma^\mu \partial_\mu - m) \psi = 0, \quad (10)$$

where $\gamma^\mu = (\beta, \beta \boldsymbol{\alpha})$ are 4×4 Dirac matrices which anticommute $\{\gamma^\mu, \gamma^\nu\} = 2g^{\mu\nu}$.

The probability amplitudes ψ, ψ^+ give 4-th component of the probability current:

$$j_\pm^0 = u_\pm^0 \rho = \pm c \cdot \psi^+ \psi, \quad (11)$$

where $u_\pm^0 = \pm c$ is the time component of the 4-speed (here $c \neq 1$).

The specifics of relativistic systems are more easily and naturally taken into account within the HJ formalism, where degrees of the Hamiltonian and 3-momentum in the evolution equation are the same. Therefore, the main relationships and equations of TS RQM will be obtained precisely in this formalism.

In the non-relativistic version of the HJ formalism, 3-momentum is defined as the gradient of the action function along trajectory $\vec{p} = -\nabla S$ and the free Hamiltonian $H_0 = (\nabla S)^2 / 2m$ enters into the HJ equation $\partial_t S + H_0 = 0$ linearly.

But (7) gives the relativistic HJ equation with a quadratic Hamiltonian:

$$(\partial_0 S)^2 = H_0^2, \quad H_0^2 = (\nabla S)^2 + m^2 \Rightarrow \partial_\mu S \cdot \partial^\mu S = m^2. \quad (12)$$

At passing to the nonrelativistic limit, the contribution of the rest energy is eliminated $S = S' - mt$, and the HJ equation for S' again becomes linear in H_0 .

In TS RQM, translational symmetry gives $\psi(x+a) = e^{-iPa} \psi(x)$ and the translation generator, 4-momentum operator $P^\mu = i\partial^\mu$, by acting on the wave function gives its eigenvalues p^μ :

$$P^\mu \psi = p^\mu \psi. \quad (13)$$

Re-acting P^μ to (13) gives, taking into account $p^\mu p_\mu = m^2$, the KG equation (8):

$$P^\mu P_\mu \psi = -\partial^\mu \partial_\mu \psi = p^\mu p_\mu \psi = m^2 \psi. \quad (14)$$

Energy expressed in terms of 3-momentum is the Hamiltonian $P_0(\vec{P}) = H_0$ and therefore:

$$-\partial_0^2 \psi = H_0^2 \psi = (\vec{P}^2 + m^2) \psi. \quad (15)$$

Here the second-order equations (14) and (15) express the relativistic relationship between energy and momentum and are therefore covariant and HJ forms of *the general wave equation*, which is satisfied by the physical components of the wave functions of all particles of finite mass, regardless of spin.

Since the beginning of the creation of RQM, adequate consideration of the specifics of relativity theory has remained only a partially solved problem, which has led to a number of misconceptions [9,10].

First, the Hamiltonian enters the general equation (14)-(15) quadratically, and this is an essentially important fact. Energy has two signs, and the state space is divided into two orthogonal parts with a specific energy sign.

Secondly, particles in these states of different energy signs go in different directions of the time axis.

Third, the wave equations in TS RQM for free particles are linear and homogeneous, and therefore the wave functions, as in the nonrelativistic case, are defined up to normalization factors. In relativity theory, these factors can change under frame-of-reference transformations; in particular, the normalization volume changes. The final normalization volume depends on the frame of reference, and only the proper normalization volume V_0 , the volume in the particle's rest frame K_0 , is invariant.

Therefore, only the normalization per unit of this volume $V_0 = 1$ is invariant, which affects the densities in this volume measured in other frames of reference. In the former RQM, the Lagrangians in the action function, integrable over the invariant 4-volume, were assumed to be invariant, ignoring the fact that the bilinear forms of the wave functions form probability densities. Ignoring the physical meaning of these bilinear forms as probability densities contradicts quantum mechanics, and when this fact is taken into account, the former Lagrangians cease to be relativistically invariant.

Fourth, the sign of the spinor normalization factor for fermions depends on the sign of energy. Therefore, the signs of the bilinear forms of the operators of physical quantities depend on the sign of the spinor normalization factor, although the eigenvalues of operators are independent of the sign of this factor.

Fifth, the relativistic HJ equation contains a squared Hamiltonian, which rules out analogy with the model of coupled harmonic oscillators, where the Hamiltonian is proportional to the square of the canonical momentum. Despite the covariant form of this model, the Hamiltonian, as a quadratic form, was effectively nonrelativistic; negative boson energy eigenvalues were excluded. Negative energies appeared only under the condition of a negative norm of states, which is not allowed by quantum mechanics. Relativity theory, however, is symmetric for states of both energy signs; in particular, in TS RQM, negative fermion energies arise naturally and are used to describe positive-energy antiparticles.

The first specificity was formerly correctly taken into account only for fermions by linearizing the square of the Hamiltonian and turning to first-order equations (13) or special wave equations. If the KG equation obeys a separate physical component of the wave function, then special wave equations determine the mixing of components due to spin and interactions.

The second specificity is taken into account according to ZSF interpretation [2-4] and TSR [7,8].

TS RQM also takes into account a third specificity of relativistic systems: the properties of the transformations of normalization volumes and densities. Since the normalization $V_0 = 1$ is invariant, the eigen wave functions ψ_0 defined in K_0 are also invariant and lead to an invariant probability density $\rho_0 = \psi_0^\dagger \psi_0$. In the laboratory frame of reference K_+ , where different particles have different 3-momenta, the corresponding to V_0 laboratory normalization volume $V(V_0)$ depends on energy of each particle: $V = m / E_p$. Accordingly, the laboratory wave functions ψ define the laboratory probability densities $\rho = \psi^\dagger \psi$. Note that in the former statement of the Lorentz invariance of the bilinear form in the mass term $\bar{\psi}\psi$ in the fermion Lagrangian, where $\bar{\psi}$ is the conjugate spinor, the fact that it is proportional to the probability density, transforming as the inverse power of the volume, was not taken into account. Therefore, in this case, the HJ formalism in its relativistic modification (12), leading to the equations (14)-(15), is more adequate than the Lagrangian formalism.

To account for the fourth specificity, we proceed from the fact that in quantum mechanics, physical quantities are given by the eigenvalues of operators F in the states ψ , which follow from the equations for eigenvalues $F\psi = f\psi$. First, these equations are homogeneous with respect to ψ , and therefore the eigenvalues f are independent of the normalization constants ψ . Second, the probability density of a physical quantity is expressed by a matrix element $\bar{\psi}F\psi = f\bar{\psi}\psi$. Here $\bar{\psi}$ is the conjugate wave function, which for bosons is equal to the Hermitian conjugation $\bar{\psi} = \psi^\dagger$, and for fermions includes the Dirac matrix: $\bar{\psi} = \psi^\dagger \gamma^0$. As we see, $\bar{\psi}F\psi$ yields not the magnitude of f , but its value multiplied by the normalization factor $\bar{\psi}\psi \neq 0$, proportional to the probability density. Therefore, the value f is given by the normalized matrix element:

$$f = \frac{\bar{\psi}F\psi}{\bar{\psi}\psi}. \quad (16)$$

In relativistic mechanics energy-momentum tensor in K_+ for a system of non-interacting particles is equal to $\varepsilon_0 u^\mu u^\nu$, i.e. it is given by the product of the invariant proper energy density $\varepsilon_0 = dm / dV_0$, given in K_0 , and two 4-velocities u^μ , specified in K_+ . In TS RQM ε_0 is given the meaning of the probability density of the rest energy in an ensemble of particles, where the probability w is given, and it differs from the intrinsic probability density $\rho_0 = dw / dV_0$ only by multiplication by the rest energy, i.e. $\varepsilon_0 = m\rho_0$. The definition of the probability density of 4-momentum, thus, takes the form $m\rho_0 u^\mu = m j^\mu$, where

$$j^\mu = \rho_0 u^\mu \quad (17)$$

is a 4-vector of the probability current density. It mu^μ is determined by the drift (real) part of the eigenvalues of 4-momentum operator $i\hbar\partial^\mu$ in K_+ , which we express through the normalized bilinear form:

$$p^\mu = \frac{\bar{\psi} i \partial^\mu \psi}{\bar{\psi} \psi}. \quad (18)$$

For a free particle, the average of the 4-current operator is then equal to the integrals over two types of volumes of this Hermitian bilinear form:

$$J^\mu = \int dV_0 \rho_0 u^\mu = \int dV \rho u^\mu = \frac{i}{2m} \int dV \sigma [\bar{\psi} \partial^\mu \psi - (\partial^\mu \bar{\psi}) \psi], \quad (19)$$

where $\sigma = \psi^\dagger \psi / \bar{\psi} \psi$. For bosons and fermions $\psi^\dagger \psi > 0$, but for negative energy of fermions $\bar{\psi} \psi < 0$ and $\sigma < 0$, and in all other cases $\bar{\psi} \psi > 0$, $\sigma > 0$.

The average of the time component of the current, as a result, has the form:

$$J^0 = \int dV_0 \rho_0 u^0 = \int dV \rho u^0 = \frac{i}{2m} \int dV \sigma [\bar{\psi} \partial_t \psi - (\partial_t \bar{\psi}) \psi]. \quad (20)$$

Since the current is negative at negative boson energies, boson theory was thought to lead to negative probability. In TS RQM, the current expression (20) is general for all relativistic particles. In the nonrelativistic limit, the probability current density takes the form (here $c \neq 1$):

$$j_\pm^0 \approx \sigma c \psi^\dagger \psi, \quad \vec{j} \approx \frac{\vec{p}}{m} \psi^\dagger \psi. \quad (21)$$

The time component $j_\pm^0$ has two signs for two directions of current along the time axis for two signs of energy when the current velocities are $(\pm c, \vec{v})$.

For free particles, the usual continuity equation is valid and σ does not depend on coordinates: $\partial_\mu j^\mu = 0$. Thus, in TS RQM j^μ is expressed through the drift part of the 4-momentum (18) with the normalization factor σ . For negative-energy fermions $\sigma < 0$, which compensates the negative sign of the form $\bar{\psi} \psi$, and $\sigma > 0$ in all other cases, which does not affect the sign.

Thus, a careful definition of normalization factors is not limited to the contribution of the volume factor and leads to a difference in sign for the two energy values for fermions. The transition to normalized averages eliminates the dependence of observable values on the sign of the spinor normalization. This leads to a correction of fermion theory, where 4-th components of probability currents and interaction currents of positive charges change sign when energy sign changes, just as in the case of bosons.

3. Time-symmetric effective quantum fields

3.1. On the transition from TS RQM to TS QFT and hypotheses about physical quantum fields

The former explanation of the transition from RQM to QFT was based on the argument that in systems with variable number of particles a theory of systems with a variable number of particles is needed, which was considered as QFT. This argument was important initially, but developments showed it to be insufficient and imprecise. The propagator approach, based on RQM and involving only *particles and external fields*, ensured multiparticle theory and led more simply to the same diagram technique that followed from QFT, incorporating the hypothesis of *physical quantum fields*.

Therefore, the question arose why for obtaining the same matrix elements sufficient to describe practically all processes in particle physics in RQM, it is necessary to switch to QFT with its additional physical quantum fields?

In the former QFT, the existence of a *physical quantum field*, filling entire space, was postulated for each type of particle [9-11]. Based on the analogy with the

electromagnetic field and photons, and the success of the second quantization method in condensed matter theory, which described phonons and other objects, QFT introduced a physical field for each fundamental particle. After quantization of the field, the particle was interpreted as an excited state of this field.

However, in condensed matter theory, it was obvious that fields were effective, and that we were talking about multiparticle systems of atoms and molecules. In QFT, quantum fields were introduced seriously as physical fields with an ontological role as primary entities, while particles were considered secondary entities.

Therefore, TS QFT returns to the questions that initially arose in connection with many hypothetical physical fields for a multitude of matter particles in SM, for which the associated external fields as sources have not been discovered:

- a) what are these hypotheses about fields based on and what confirms their consequences;
- b) what to do with all the problems arising from these hypotheses;
- c) what is the inadequacy of RQM formalism, which does without such fields?

In reality, hypotheses about quantum fields in addition to known external fields lacked experimental support and led to insoluble problems. One of the main and directly verifiable consequences of quantum fields was the zero-point fluctuations of their vacuums, which led to an infinite zero-point energy of the vacuums of all such fields in the vicinity of each point. This made QFT incompatible with gravity, which does not allow for a simple shift in energy levels.

For the electromagnetic field, QFT prediction of zero-point fluctuations in the quantum field vacuum was *rejected* by experiments shortly after the creation of covariant RQM and QFT. This obvious fact, however, has since been interpreted in the opposite way, as *confirmation* of the existence of zero-point fluctuations.

First, measurements of the Lamb shift and anomalous magnetic moments revealed only the contribution of loop corrections from RQM, i.e., the contribution of electron-photon interactions. The contribution of zero-point vacuum fluctuations from QFT, which would have doubled the effect, was completely absent in the experiments.

Secondly, measurements of the Casimir effect also showed that there are only contributions from external fields of real sources, the atoms of the crystal, and there is no contribution from zero-point oscillations of the quantum field vacuum, which would double the effect here too.

In QFT, the two mechanisms that should contribute to each of these effects - the contributions of particles in loop diagrams and the contributions of vacuum fluctuations - were considered to be two sides of the same phenomenon. However, these assumptions were initially incorrect for both experiments. In the first case, the loop contributions cannot be neglected, while in the second, fluctuating external fields generated by atoms in crystals must be taken into account. It turned out that these real contributions are sufficient to describe the entire observed effect in both cases.

If we add to them the contributions of hypothetical zero-point fluctuations of the quantum field vacuum, then the QFT predictions for these effects would become twice as large as those observed experimentally. And this is at conditions where measurement errors are much smaller than the effect. Such a catastrophic discrepancy between theory and experiment seemed impossible given the remarkable successes of QFT.

In reality, all the successes of QFT were related to the same methods already used in RQM – solution of relativistic equations for wave functions and diagram techniques. Therefore, ignoring the obvious discrepancies between the former QFT and experiment was due to a reluctance to accept the fact that hypotheses about physical quantum fields had been rejected by experiments.

In this regard, questions about the relationship between earlier forms of RQM and QFT are of purely historical significance and persist in the literature partly due to the

inertia of thought and partly due to entrenched dogmatism regarding historical legacy. This situation is analogous to the history of the ether in pre-relativistic physics. The analogy with waves in continuous media gave rise to the hypothesis about the ether as a medium for electromagnetic waves. Experiments and relativity theory rejected this hypothesis, yet it remained for a long time in scientific literature and education system.

Now it turns out that the external electromagnetic field requires neither ether nor the addition of a physical quantum field. In TS QFT, the quantum, operator part of the field is formally added, but as an effective field. Being the operator amplitude of a transition in the space of occupation numbers, it transforms the states of a multiparticle system. In this regard, we further notice the following.

First, the quantum-mechanical harmonic oscillator can be described both in terms of wave functions and their corresponding energy levels and in terms of raising-lowering operators with the same energy levels, and both approaches are equivalent. The descriptions of *free* relativistic particles in RQM and QFT are analogous to these approaches to the oscillator and are therefore also equivalent.

Secondly, there is a well-known classical field, the electromagnetic field, whose quantum properties boil down to the fact that this field, while remaining an external field, can emit and absorb photons, the quanta of this field, which then behave like particles. That is, local quantum perturbations of this external field arise, move, and are absorbed like particles. However, this does not allow one to attribute continuous quantum fluctuations to the field throughout space. Therefore, for external fields, quantum properties are relevant only locally, at the points of emission and absorption of quanta. Between these events, we are dealing with particles and external fields only, when RQM and QFT are essentially equivalent.

Third, even describing particle emission and absorption processes as quanta of interaction fields does not mean that RQM is no longer applicable here and that QFT cannot be dispensed with. At expanding the *S*-matrix into a perturbation theory series with sources and the external fields, these sources generate yields the same matrix elements for the processes in RQM as QFT with its field operators.

For these reasons, at describing both free and interacting particles, there are no fundamental reasons to abandon RQM and move to QFT, since both theories, after discarding the “extra” part of QFT, the quantum field hypothesis with its physical vacuum, ultimately lead to the same transition amplitudes.

However, it's a completely different matter when QFT (with its fitting procedures) is considered not a more correct theory, replacing the flawed RQM, but an alternative approach with a more effective formalism. In practice, the transition from RQM to QFT occurs precisely this way, since, aside from initial differences, the results of RQM and QFT are identical, and the problematic parts of QFT are simply ignored.

Here, we'll pose the question differently and ask whether it's possible to formulate QFT as a consistent and fundamental theory, similar to RQM, without the problems that plagued the former QFT. To answer this question, we'll consider two key facts.

Firstly, as has been noted above, RQM with external fields and particles allows one to describe almost all processes of particle physics [5-7].

Secondly, there are no grounds for hypotheses about quantum fields, and all purely quantum field problems of QFT arise precisely from them.

Therefore, TS QFT will be constructed further without these hypotheses, restricting itself, as in TS RQM, *to external fields and particles*. As will be shown below, this becomes natural once we consider the fact that the field operators are, in fact, the former transition probability amplitudes in a multiparticle system. The transition to QFT methods will therefore not mean the addition of physical quantum fields, but what it has always essentially been - a transition to the theory of multiparticle systems with effective field functions and operators.

3.2. Quantum field as an effective field for describing a multiparticle system

As noted above, the theories of boson fields in the former QFT were constructed by analogy with a chain of coupled harmonic oscillators. The models had an apparently covariant form, but the Hamiltonian was proportional to the square of the canonical momentum, included the zero-point vacuum energy, and, as a quadratic form, had only positive eigenvalues [9-11]. All of this is characteristic of a nonrelativistic system.

The positive energy made the model incompatible with ZSF interpretation, which has no ground-state zero-point energy. Introducing negative energy into bosonic field models led to negative norms for such states. Thus, the models with antiparticles led to a divergent ground-state energy, while in ZSF interpretation the norms became negative. This rendered the bosonic field theory of the former QFT physically inconsistent.

Thus, former analogies with classical models were misleading and led to insoluble problems. Therefore, we will now examine the introduction of effective field operators, taking into account the specifics of relativistic and quantum theories for systems with a variable number of particles.

First, relativistic covariance requires the same degrees of energy and momentum for particles. In general, this led to a quadratic HJ equation for single-particle wave functions and field operators. This general equation is sufficient for bosons, but for fermions, where spin mixes components, a linearization of the HJ equation is required.

Second, in the second-quantization method, field operators change the number of particles in states and are formed from single-particle wave functions and operators that increase or decrease the number of particles in the system. They are similar to quantum field operators, but as probability amplitudes, they satisfy wave equations.

Third, one of the reasons for the problems of the former QFT was the introduction of Lagrangians and Hamiltonians for quantum fields in addition to the particle Hamiltonians.

Taking these circumstances into account radically changes the physical meaning and the very concept of the quantum field theory, simplifying and eliminating artificially created problems. In TS QFT, the Hamiltonians of particle systems are formed from the Hamiltonians of the particles, adding only parts with external fields and interactions of the particles themselves. In this approach, the quantum field is not a new entity in addition to particles, as formerly believed, but an effective field introduced to conveniently describe a multiparticle system.

While the former QFT extended the quantization of several interaction fields to all known particles, TS QFT is based on the opposite approach: the practice of describing multiparticle systems of particles, for which only probability amplitudes are specified in space, also extends to the quanta of interaction fields. As physical fields, however, remain only external fields, acting as sources for their quanta. Unlike the quanta of interaction fields, other particles of matter are not assigned physical fields as their sources. Therefore, they are either created and disappear as particle-antiparticle pairs, or their mutual transformations are introduced on the basis of particle physics models. The ground state of a multiparticle system is therefore a state in which there are no particles, in contrast with the vacuum of quantum fields with divergent energy density.

As a result, QFT not only resolves the problems of particle physics but also challenges it to create and test models of the mutual transformations of particles still considered as fundamental ones. This situation is similar to the postulation of "quantum fields" for each hadron before the composite nature of hadrons was elucidated, but this practice was later replaced by mechanisms of local transformations of some particles into others through the redistribution of quarks.

For now, let us begin with the consequences of the fact that in TS QFT, wave functions $\psi(x)$ become formally field operators, without yet clarifying the physical meaning of such a change and being interested only in its general consequences [9-11].

4-momentum operator P_μ is a generator of translations in spacetime and gives the dependence of operators on coordinates: $\psi(x) = e^{iP_\mu x} \psi(0) e^{-iP_\mu x}$. The derivatives $i\partial_\mu$ are equal to commutators P_μ with the field operator ψ :

$$i\partial_\mu \psi = [\psi, P_\mu], \quad i\partial_\mu \bar{\psi} = [\bar{\psi}, P_\mu]. \quad (22)$$

As shows (22), the eigenvalues of the bilinear form:

$$i\bar{\psi} \partial_\mu \psi = \bar{\psi} [\psi, P_\mu] = \bar{\psi} \psi P_\mu - P_\mu \bar{\psi} \psi \quad (23)$$

should be interpreted, as in TS RQM (19), as the operator of 4-momentum of a state. From this follow the expressions for the Hamiltonian and 3-momentum:

$$H_0 = \int d^3x \bar{\psi} i\partial_0 \psi, \quad (24)$$

$$\mathbf{P} = -\int d^3x \bar{\psi} i\nabla \psi. \quad (25)$$

Re-applying $i\partial_\mu$ and (22) taking into account the conservation of energy lead to:

$$-\partial_\mu \partial^\mu \psi = [i\partial^\mu \psi, P_\mu] = [[\psi, P^\mu], P_\mu]. \quad (26)$$

At acting on a single-particle state Φ_p , three of the four terms on the right (26) give zero $P^\mu \psi \Phi_p \sim P^\mu \Phi_0 = 0$, since the field operator transforms the single-particle state into the ground state $\psi \Phi_p \sim \Phi_0$. The remaining term gives $\psi P^\mu P_\mu \Phi_p = m^2 \psi \Phi_p$ and from (26) it follows the KG equation for the field operator ψ :

$$(\partial_\mu \partial^\mu + m^2) \psi = 0. \quad (27)$$

The same equation is also valid for a multiparticle state, since the operator $\psi(x_i)$ acts each time on the state of one of the particles with coordinate x_i .

The KG equation (27), as in the TS RQM, is a differential form of the relativistic HJ equation:

$$-\partial_0^2 \psi = H_0^2 \psi, \quad H_0^2 = (\mathbf{P}^2 + m^2). \quad (28)$$

The equations (27)- (28), expressing the relativistic relationship between energy and momentum of particles, are two (differential and Hamiltonian) forms of the *general field equation* for the component of any field operator for the system of relativistic particles.

In the case of massless particles, the differential and Hamiltonian forms of the HY equation (27)-(28) are simplified and take the form:

$$\partial_\mu \partial^\mu \psi = 0, \quad (29)$$

$$-\partial_0^2 \psi = H_0^2 \psi, \quad H_0^2 = \mathbf{P}^2. \quad (30)$$

In a stationary state Φ_p the eigenvalue equation:

$$H_{(0)}^2 \psi \Phi_p = E_p^2 \psi \Phi_p, \quad (31)$$

yield $E_p^2 = \mathbf{p}^2 + m^2$, where $\mathbf{p}$ is the eigenvalue of $\mathbf{P}$. From (24) and (31) it follows that the eigenvalues $\pm E_p$ of the Hamiltonian H_0 in states of two energy signs with field operators $\psi_{\pm}$ are given by the equations:

$$H_0 \psi_{\pm} \Phi_p = \pm E_p \psi_{\pm} \Phi_p. \quad (32)$$

Thus, the solutions of the KG equation (27), the general field equation for the physical components of free relativistic fields, are complex field operators $\psi_{\pm}(\mathbf{x}, t)$ acting in two state spaces with energies $p_{0\pm} = \pm E_p$. These states $\Phi_{\pm p}$ evolve, depending on the sign of energy, in two directions of times t and are normalized in their spaces.

Equations (22)-(32) allow one to determine the physical meaning of a field operator that has two components with different energy signs $\psi(x) = \psi_+(x) + \psi_-(x)$. Let us consider the momentum decomposition of these components, taking into account mass-shell constraint:

$$\psi_+(x) = \int d^4 p e^{-ipx} \theta(p^0) \delta(p^2 - m^2) \psi_+(p), \quad (33)$$

$$\psi_-(x) = \int d^4 p e^{-ipx} \theta(-p^0) \delta(p^2 - m^2) \psi_-(p). \quad (34)$$

Substituting this into (22) gives:

$$[\psi_+(x), P^{\mu}] = \int d^4 p e^{-ipx} p^{\mu} \theta(p^0) \delta(p^2 - m^2) \psi_+(p), \quad (35)$$

$$[\psi_-(x), P^{\mu}] = - \int d^4 p e^{-ipx} p^{\mu} \theta(-p^0) \delta(p^2 - m^2) \psi_-(p), \quad (36)$$

from which it follows:

$$[\psi_+(p), P^{\mu}] = p^{\mu} \psi_+(p), \quad (37)$$

$$[\psi_-(p), P^{\mu}] = -p^{\mu} \psi_-(p). \quad (38)$$

Considering:

$$P^{\mu} \Phi_{\pm p'} = p'^{\mu} \Phi_{\pm p'}, \quad (39)$$

we multiply (37)-(38) on the right by $\Phi_{\pm p'}$ and get:

$$P^{\mu} \psi_{\pm}(p) \Phi_{\pm p'} = (p'^{\mu} \mp p^{\mu}) \psi_{\pm}(p) \Phi_{\pm p'}, \quad (40)$$

$$P^{\mu} \psi_{\pm}^+(p) \Phi_{\pm p'} = (p'^{\mu} \pm p^{\mu}) \psi_{\pm}^+(p) \Phi_{\pm p'}. \quad (41)$$

It follows that $\psi_{\pm}(p) \Phi_{\pm p'}$ is an amplitude for a state with energy-momentum $p' \mp p$, and $\psi_{\pm}^+(p) \Phi_{\pm p'}$ is an amplitude for a state with energy-momentum $p' \pm p$. Given the relation $p^2 = m^2$, we conclude that the operator $\psi_{\pm}^+(p)$ adds a particle of mass m and 4-momentum $\pm p$, while the operator $\psi_{\pm}(p)$ subtracts such a particle.

Thus, the physical meaning of effective field operators $\psi_{\pm}(x)$, $\psi_{\pm}^+(x)$ is that they are operator amplitudes for the transition of a multiparticle system into states that differ by one particle.

Therefore, field operators expanded in free wave functions lead to an equidistant spectrum of free particle states. Multiparticle states of non-interacting particles are

formed as products of single-particle states, and the procedures for their formation described for the former QFT also applicable to TS QFT, with the corresponding replacement of antiparticle states with negative-energy states.

3.3. Time-symmetric quantization of boson fields

Field operators for states of a given energy sign are complex, since they represent wave functions multiplied by operators increasing or decreasing the number of particles in the system. Even combining field components of two energy signs does not lead to a real field operator, since only operators with the same energy sign are mutually Hermitian adjoint. Therefore, in TS QFT, all field operators are complex, regardless of the presence or absence of spin and charge.

In the spinless case, the effective field operator has one component for each energy sign. The vector field has three physical components, also for each energy sign. For particles of the massless gauge field (photon and gluon fields) and the graviton field, there are two transverse components, forming a single complex field operator (see Section 3.5).

To obtain solutions for the physical components of the free bosonic field operators required for SM the general KG equation (27) and the general formula for the Hamiltonian (24) are sufficient. Accounting for polarizations and separating transverse components in different reference frames for vector and tensor fields are the technical problem solved by introducing appropriate projectors. However, this complication has virtually no effect on the quantization of the physical components thus identified. Therefore, it is sufficient to describe the quantization of the effective fields introduced for spinless particles.

The expansions of the complex field operator in solutions of the general field equation, i.e. the KG equation (27), for two signs of energy have the form:

$$\phi(x) = \phi_-(x) + \phi_+(x) = \int d\tilde{p} (a_p e^{-ipx} + a_{-p} e^{ipx}), \quad (42)$$

$$\phi^+(x) = \phi_+^+(x) + \phi_-^+(x) = \int d\tilde{p} (a_{+p}^+ e^{ipx} + a_{-p}^+ e^{-ipx}), \quad (43)$$

where $d\tilde{p} = d^3\mathbf{p} / (2\pi)^3 2E_p$, and $a_{\pm p}^+, a_{\pm p}$ are of creation and annihilation operators.

In (42)-(43) it is clearly seen that the field operators ϕ, ϕ^+ , expanded in terms of single-particle wave functions, have the same meaning as probability amplitudes as the wave functions, but when multiplied by operators that increase or decrease the number of particles in the system, they become operators expressing transition probabilities in the space of occupation numbers. Therefore, like wave functions in quantum mechanics, they form a canonical pair with commutators:

$$[\phi(\mathbf{x}, t), \phi^+(\mathbf{x}', t)] = \delta^3(\mathbf{x} - \mathbf{x}'), \quad (44)$$

$$[\phi(\mathbf{x}, t), \phi(\mathbf{x}', t)] = [\phi^+(\mathbf{x}, t), \phi^+(\mathbf{x}', t)] = 0. \quad (45)$$

These commutators are similar to fermionic anticommutators in that they also do not contain time derivatives. From (42)-(45) we obtain commutators of the creation-annihilation operators of quanta of two energy signs:

$$[a_{\pm p}, a_{\pm p'}^+] = 2E_p^2 (2\pi)^3 \delta^3(\mathbf{p} - \mathbf{p}'), \quad (46)$$

$$[a_{\pm p}, a_{\pm p'}] = [a_{\pm p}^+, a_{\pm p'}^+] = 0. \quad (47)$$

Two important differences from standard QFT are noted here: the commutators are positive-definite for both energy signs, and the state energy enters them quadratically E_p^2 . The quadratic dependence on energy in (46) is also similar to the case of fermions,

where the two energy degrees are distributed to the anticommutator of the creation-annihilation operators and the bilinear form of the bispinors.

The annihilation operators define the ground states of two energy signs:

$$a_{+p}|0\rangle = 0, \quad a_{-p}|0\rangle = 0, \quad (48)$$

and the creation operators determine the corresponding single-particle states:

$$a_{+p}^+|0_+\rangle = |p\rangle, \quad a_{-p}^+|0_-\rangle = |-p\rangle. \quad (49)$$

The Hamiltonian of the field (24) and momentum (25) in terms of $a_{\pm p}^+$ and $a_{\pm p}$ have the form:

$$H = \int d\tilde{p} (a_{+p}^+ a_{+p} - a_{-p}^+ a_{-p}), \quad (50)$$

$$\mathbf{P} = \int d\tilde{p} \frac{\mathbf{P}}{E_p} (a_{+p}^+ a_{+p} - a_{-p}^+ a_{-p}). \quad (51)$$

As can be seen from (50), the ground state energy is zero, i.e., there is no zero-point energy in TS QFT. The particle number operator $N = \phi^+ \phi$ is:

$$N = \int \frac{d\tilde{p}}{E_p} (a_{+p}^+ a_{+p} + a_{-p}^+ a_{-p}). \quad (52)$$

Thus, in boson theory, as in TS RQM, the state norms are positive, the probability current is formed from 4-momentum of the state and there is no zero-point energy in the ground state.

3.4. Time-symmetric quantization of the fermion field of spin $\frac{1}{2}$

The specificity of relativity theory - the quadratic Hamiltonian at quadratic 3-momentum - was correctly taken into account in RQM only for spin $\frac{1}{2}$ fermions. We choose the matrix Hamiltonian, obtained by linearizing the quadratic relativistic expression (28), in its minimal form:

$$H_0 = \psi^+ (-i\boldsymbol{\alpha} \cdot \nabla + \beta m) \psi. \quad (53)$$

Here the field equations are the Dirac equation, a *special field equation* that determines the mixing of field components due to spin:

$$(i\gamma^\mu \partial_\mu - m)\psi = 0. \quad (54)$$

The Hamiltonian (53), taking into account (54), takes the form:

$$H_0 = \psi^+ i\partial_0 \psi. \quad (55)$$

The momentum expansion of the field operators has the form:

$$\psi(x) = 2m \sum_{\alpha=1}^2 \int d\tilde{p} (b_{p\alpha} u_p^\alpha e^{-ipx} + b_{-p\alpha} v_p^\alpha e^{ipx}), \quad (56)$$

$$\psi^+(x) = 2m \sum_{\alpha=1}^2 \int d\tilde{p} (b_{p\alpha}^+ u_p^{\alpha+} e^{-ipx} + b_{-p\alpha}^+ v_p^{\alpha+} e^{ipx}). \quad (57)$$

Since fermion wave functions are antisymmetric, field operators must change sign at rearrangement. Therefore, when the operator transition amplitudes in commutators are rearranged, their product changes sign, and the commutators transform into anticommutators, the nonzero of which has the form:

$$\{\psi^\alpha(\mathbf{x}, t), \psi^{+\alpha'}(\mathbf{x}', t)\} = \delta^3(\mathbf{x} - \mathbf{x}') \delta^{\alpha\alpha'} I. \quad (58)$$

Substitution (56)-(57) into (58), and taking into account the normalization of spinors:

$$\bar{u}_p^\alpha u_p^{\alpha'} = \delta^{\alpha\alpha'}, \quad \bar{v}_p^\alpha v_p^{\alpha'} = -\delta^{\alpha\alpha'}, \quad \bar{u}_p^\alpha v_p^{\alpha'} = 0, \quad u_p^{\alpha+} u_p^{\alpha'} = v_p^{\alpha+} v_p^{\alpha'} = \frac{E_p}{m} \delta^{\alpha\alpha'}. \quad (59)$$

gives an expression for the anticommutators:

$$\{b_{\pm p\alpha}, b_{\pm p'\alpha'}^+\} = (E_p / m)(2\pi)^3 \delta^3(\mathbf{p} - \mathbf{p}') \delta_{\alpha\alpha'}. \quad (60)$$

The annihilation operators define the ground states for both types of particles:

$$b_{p\alpha} |0_+\rangle = 0, \quad b_{-p\alpha} |0_-\rangle = 0, \quad (61)$$

and the creation operators determine the corresponding single-particle states:

$$b_{p\alpha}^+ |0_+\rangle = |p\rangle, \quad b_{-p\alpha}^+ |0_-\rangle = |-p\rangle. \quad (62)$$

The expression for the Hamiltonian (55) and the number of particles in terms of the creation-annihilation operators take the form:

$$H_0 = 2m \sum_\alpha \int d\tilde{p} E_p (b_{p\alpha}^+ b_{p\alpha} - b_{-p\alpha}^+ b_{-p\alpha}), \quad (63)$$

$$N = 2m \sum_\alpha \int d\tilde{p} (b_{p\alpha}^+ b_{p\alpha} + b_{-p\alpha}^+ b_{-p\alpha}). \quad (64)$$

Thus, there is no zero-point energy for fermions in the ground state either.

3.5. Massless SM boson fields and gravitons

Let's consider the physical aspects of describing massless quanta emitted or absorbed by external gauge fields and an external gravitational field. Effective field operators here also include operator transition amplitudes, which determine the form of their commutators.

For our purposes, it is sufficient to deal with the physical components of these field operators in the source rest frame, and then the results can be transformed into other reference frames.

Let us consider the quanta of transverse waves of the circularly polarized electromagnetic field, which propagate as photons, massless particles of spin 1. Under the conditions $\partial_\mu A^\mu = 0$, $\mathbf{k} = (0, 0, k^3)$ and gauge $A_3 = A_0 = 0$, transverse components $A = (A_1 + iA_2)/\sqrt{2}$, $A^+ = (A_1 - iA_2)/\sqrt{2}$ are similar to the operator of a complex spinless field with a chiral charge, the role of which is played by helicity Λ .

In TS QFT, these two states with opposite helicities, acting as a particle-antiparticle pair, are described as states of two energy signs $A_\pm$. The free Hamiltonian and the number operator are therefore given by expressions as for the spinless field:

$$H_0 = \int d^3x A^+ i\partial_t A, \quad (65)$$

$$N = \int d^3x A^+ A. \quad (66)$$

The momentum expansion of the corresponding solution of the field equation is similar (56)-(57) with creation-annihilation operators $a_{\pm k}^+, a_{\pm k}$ for states with the same helicity $\Lambda_\pm = +1$, but with different energy signs.

For arbitrary orientation of the coordinate axes, a projection operator will emerge that isolates the transverse components. The state probabilities remain positive-definite, similar to the case of a spinless field. The observable operators are also similar to (63)-(64), from which it is clear that the ground states do not contain zero-point energy.

In the case of circularly polarized wave quanta of a massless non-Abelian vector gauge field A_μ^a (spin 1) and transverse gravitational wave quanta (spin 2), there are also two physical degrees of freedom and axial symmetry. Therefore, their quantum theory is similar to the case of photons, and there is no zero-point energy in the ground state. Further consideration of nonlinearity, i.e., interactions between field components, and internal symmetries does not change the conclusion that there is no zero-point energy in the free Hamiltonian.

The conclusion that for massless gauge fields and the gravitational field in one frame of reference and in the transverse gauge there is no zero-point energy remains valid for all frames of reference and gauges due to the invariance of the ground state.

The field operators of the two transverse components of the gauge field $W^\pm$ are mutually charge-conjugate, and before they are given mass by the background scalar field, their quanta behave as particle-antiparticle. In the case of the components of the neutral gauge field Z_L, Z_R and two neutral scalar fields, these quanta, before they are given mass by the interaction terms, were also mutually charge-conjugate. There was no zero-point energy in the free Hamiltonians of $W^\pm$ and Z_L, Z_R , and the inclusion of interactions with the scalar field (to impart an effective mass) does not generate a zero-point energy for the ground state of the free field.

From a more general perspective, in TS QFT, there is no particle without an antiparticle, since the relativistic field is described by a complex field due to the presence of two energy signs. The similarity of the currently known quantum numbers of a particle and its antiparticle merely leads to their practical indistinguishability, but this does not mean their complete identity. As a rule, there is at least a very weak interaction, which leads to their distinguishability.

4. Commutators, time-symmetric chronological ordering, and causal propagators

4.1. Commutator of fields and microcausality

Commutators of spinless field operators in the general case are equal to:

$$\begin{aligned} [\phi(\mathbf{x}', t'), \phi^+(\mathbf{x}, t)] &= \int \frac{d^3\mathbf{k}}{(2\pi)^3 2} e^{i\mathbf{k}(\mathbf{x}'-\mathbf{x})} (e^{-iE_k(t'-t)} + e^{iE_k(t'-t)}) = \\ &= -i\partial_t \int \frac{d^3\mathbf{k}}{(2\pi)^3 2E_k} e^{i\mathbf{k}(\mathbf{x}'-\mathbf{x})} (e^{-iE_k(t'-t)} - e^{iE_k(t'-t)}) = \partial_t D(\mathbf{x}'-\mathbf{x}), \end{aligned} \quad (67)$$

where $D = D^+ + D^-$ - the Pauli-Jordan function, disappearing behind the light cone:

$$D(x) = -i \int \frac{d^3\mathbf{k}}{(2\pi)^3 2E_k} (e^{-ikx} - e^{ikx}) = -i \int \frac{d^3\mathbf{k} e^{i\mathbf{k}x}}{(2\pi)^3 2E_k} (e^{-iE_k t} - e^{iE_k t}). \quad (68)$$

This is consistent with simultaneous commutators (44).

The bilinear form of field operators is proportional to the probability density, which transforms as the time component of a vector, and the equality of the commutator function in (67) to the time derivative of the invariant function $D(x)$ clearly demonstrates this fact. Introducing a normal $n^\mu = (1, 0, 0, 0)$ to the hypersurface of simultaneity, where the probabilities are given, we obtain:

$$[\phi(x'), \phi^+(x)] = \partial_t D(x' - x) = n^\mu \partial_\mu D(x' - x). \quad (69)$$

For a fermion field, the commutator of a canonical pair turns into an anticommutator and the result has the standard form:

$$\{\psi(x'), \bar{\psi}(x)\} = -(i\gamma^\mu \partial_\mu + m) iD(x' - x). \quad (70)$$

This is also consistent with simultaneous anticommutators (58).

Thus, unlike the former QFT, where there was a fundamental difference between bosonic and fermionic (anti)commutators, in TS QFT they are similar and in both cases are expressed through derivatives of the function $D(x)$. Therefore, the (anti)commutators of field operators also disappear beyond the light cone.

This fact was formerly considered as an expression of the microcausality property, a relativistic constraint that space-like events cannot be causally related [9-11]. In fact, microcausality in such a strict form is characteristic of classical (not quantum) physics, and in quantum theory is a constraint only for single-particle theory.

In relativistic quantum theory, the causal connection between two events is generally more appropriately associated with causal propagators, which will be discussed below in Section 4.3. These are expressed through a causal function $D_c(x' - x)$, which does not vanish beyond the light cone but declines rapidly beyond the Compton length of the particles. Such a correlation does not violate relativistic causality, since it is associated with the many-particle nature of processes, where the production of particle-antiparticle pairs effectively yields a formally nonlocal correlation.

4.2. Time-symmetrical chronological ordering

Chronological ordering operator T_+ of the former QFT acts only on operators of positive energy states, arranging them in order of *increasing* time from right to left:

$$T_+[A_+(t')B_+(t)] = \begin{cases} A_+(t')B_+(t), & t' > t, \\ B_+(t)A_+(t'), & t' < t. \end{cases} \quad (71)$$

In the TS QFT, the inverse to T_+ operator is also introduced, antichronological ordering operator T_- , acting only on the operators of negative energy states by ordering them in order of *decreasing* time from right to left:

$$T_-[A_-(t')B_-(t)] = \begin{cases} B_-(t)A_-(t'), & t' > t, \\ A_-(t')B_-(t), & t' < t. \end{cases} \quad (72)$$

In the general case, a time-symmetric ordering operator is introduced $\hat{T}$, which is the product of these two operators $\hat{T} = T_+T_- = T_-T_+$. For each class of operator transition amplitudes, it reduces to their time ordering operators: $\hat{T}A_+ = T_+A_+$, $\hat{T}A_- = T_-A_-$, acting selectively on the operators of both energy signs and ordering them in mutually inverse order. In particular, for, $t' > t$ we obtain:

$$\hat{T}[A_+(t')B_+(t)A_-(t')B_-(t)] = [A_+(t')B_+(t)][B_-(t)A_-(t')]. \quad (73)$$

4.3. Causal propagators of fields

In the second paper [8], time-symmetric propagators were defined as Green's functions for standard wave equations with a source $j(x)$ under the ZSF boundary conditions. They coincide with the causal propagators used in the covariant diagram technique of the propagator approach.

In TS QFT, causal propagators for bosons $G_c(x - x')$ and fermions $S_c(x - x')$ follow from the time-symmetric chronological products of the operator transition amplitudes.

Causal propagators for a complex scalar field $G_{c\pm}$ follow from the products of the operator transition amplitudes ordered in two time directions:

$$iD_{c\pm}(x' - x) = \langle 0 | T_{\pm} [\phi_{\pm}(x') \phi_{\pm}^{\dagger}(x)] | 0 \rangle. \quad (74)$$

Substituting the frequency decompositions yields:

$$\begin{aligned} iD_{c+}(x' - x) &= \langle 0 | T_{+} [\phi_{+}(x') \phi_{+}^{\dagger}(x)] | 0 \rangle = \\ &= \langle 0 | \phi_{+}(x') \phi_{+}^{\dagger}(x) | 0 \rangle \theta(t' - t) + \langle 0 | \phi_{+}^{\dagger}(x) \phi_{+}(x') | 0 \rangle \theta(t - t') = \\ &= \int d\tilde{k} d\tilde{k}' [\langle 0 | a_{k'} a_k^{\dagger} | 0 \rangle \theta(t' - t) + \langle 0 | a_k^{\dagger} a_{k'} | 0 \rangle \theta(t - t')] e^{-i(k'x' - kx)} = \\ &= \int d\tilde{k} \theta(t' - t) e^{-ik(x' - x)}. \end{aligned} \quad (75)$$

Here the matrix element $\langle 0 | \phi_{+}^{\dagger}(x) \phi_{+}(x') | 0 \rangle$ disappears and only the contribution remains $\langle 0 | \phi_{+}(x') \phi_{+}^{\dagger}(x) | 0 \rangle$. Similarly, for negative energy:

$$\begin{aligned} iD_{c-}(x' - x) &= \langle 0 | T_{-} [\phi_{-}(x') \phi_{-}^{\dagger}(x)] | 0 \rangle = \\ &= \langle 0 | \phi_{-}^{\dagger}(x) \phi_{-}(x') | 0 \rangle \theta(t' - t) + \langle 0 | \phi_{-}(x') \phi_{-}^{\dagger}(x) | 0 \rangle \theta(t - t') = \\ &= \int d\tilde{k} d\tilde{k}' [\langle 0 | a_{-k}^{\dagger} a_{-k'} | 0 \rangle \theta(t' - t) + \langle 0 | a_{-k'} a_{-k}^{\dagger} | 0 \rangle \theta(t - t')] e^{i(k'x' - kx)} = \\ &= - \int d\tilde{k} \theta(t - t') e^{ik(x' - x)}. \end{aligned} \quad (76)$$

The full causal propagator therefore takes the standard form:

$$\begin{aligned} iD_c(x' - x) &= \langle 0 | \hat{T} \phi(x') \phi^{\dagger}(x) | 0 \rangle \\ &= \langle 0 | T_{+} [\phi_{+}(x') \phi_{+}^{\dagger}(x)] + T_{-} [\phi_{-}(x') \phi_{-}^{\dagger}(x)] | 0 \rangle = \\ &= \langle 0 | \phi_{+}(x') \phi_{+}^{\dagger}(x) | 0 \rangle \theta(t' - t) + \langle 0 | \phi_{-}(x') \phi_{-}^{\dagger}(x) | 0 \rangle \theta(t - t') = \\ &= \int d\tilde{k} [\theta(t' - t) e^{-ik(x' - x)} + \theta(t - t') e^{ik(x' - x)}]. \end{aligned} \quad (77)$$

The standard result (77) also leads to the standard momentum representation:

$$iD_c(x' - x) = \frac{1}{(2\pi)^4} \int d^4k \frac{e^{-ik(x' - x)}}{k^2 - m^2 + i\varepsilon}. \quad (78)$$

Causal propagators for fermions $S_{c\pm}(x)$ also follow from the products of the operator transition amplitudes ordered in two time directions:

$$\begin{aligned}
iS_c(x'-x) &= \langle 0|T_+[\psi_+(x')\bar{\psi}_+(x)]|0\rangle + \langle 0|T_-[\psi_-(x')\bar{\psi}_-(x)]|0\rangle = \\
&= \langle 0|\psi_+(x')\bar{\psi}_+(x)|0\rangle\theta(t'-t) + \langle 0|\psi_-(x')\bar{\psi}_-(x)|0\rangle\theta(t-t') = \\
&= (2m)^2 \sum_{\alpha'\alpha} \int d\tilde{p} d\tilde{p}' [\theta(t'-t) \langle 0|b_{p'\alpha} b_{p\alpha}^+|0\rangle u_{p'}^{\alpha'} \bar{u}_p^\alpha e^{-i(p'x'-px)} + \\
&\quad + \theta(t-t') \langle 0|b_{-p'\alpha'} b_{-p\alpha}^+|0\rangle v_{p'}^{\alpha'} \bar{v}_p^\alpha e^{i(p'x'-px)}],
\end{aligned} \tag{79}$$

and finally:

$$\begin{aligned}
iS_c(x'-x) &= \int d\tilde{p} \sum_p e^{ip(x'-x)} \times \\
&\times \left[\theta(t'-t)(\gamma p + m) e^{-i\omega_p(t'-t)} + \theta(t-t')(-\gamma p + m) e^{i\omega_p(t'-t)} \right] = \\
&= -(i\gamma \partial + m) D_c(x'-x).
\end{aligned} \tag{80}$$

Thus, the ordering of the operator transition amplitudes by the time-symmetric operator $\hat{T}$ leads to the standard causal propagators.

5. Time-symmetric theory of interacting fields

5.1. Interacting particles and fields

The transition from free particles to interacting ones in TS QFT does not change the physical meaning of the effective field operators. They remain operator amplitudes for the transition between states of a multiparticle system, differing in the number and quantum numbers of particles. Interactions only change their evolution and link different multiparticle states.

The total Hamiltonian of the system is represented as $H = H_0 + H_I$, where H_0 is the Hamiltonian of free particles, and H_I describes the interactions. If an external physical field is present, its energy and interaction with the particles are also included in the total Hamiltonian of the system.

Unlike the former QFT, no additional physical Hamiltonian of certain quantum fields filling all space is introduced here. The field operators $\psi(x)$, $\phi(x)$ etc. are effective operator transition amplitudes, so the interaction Hamiltonian is constructed from known external fields and transition operators between states of the multiparticle system.

The relativistic form of the interaction must be consistent with the general HJ equation. With minimal interaction with the gauge field, the generalized derivatives and the relationship between energy and momentum take the form

$$P_\mu \rightarrow \Pi_\mu = P_\mu - qA_\mu, \quad \Pi_\mu \Pi^\mu = m^2, \tag{81}$$

or

$$(H - qA_0)^2 = m^2 + (\mathbf{P} - q\mathbf{A})^2, \tag{82}$$

In Hamiltonian form it has two energy sections

$$H_\pm = qA_0 \pm \sqrt{m^2 + (\mathbf{P} - q\mathbf{A})^2}. \tag{83}$$

The sign before the root corresponds to two signs of energy for two directions of evolution along the time axis.

The interaction Hamiltonian is determined by the difference $H_I = H - H_0$. For bosons, the expansion of the root in (83) will lead to the series of $q\mathbf{A}$:

$$\begin{aligned} H_I &= qA_0 + qV_1[\mathbf{A}] + q^2V_2[\mathbf{A}^2] + q^3V_3[\mathbf{A}^3] + \dots = \\ &= qA_0 + \sum_{n \geq 1} q^n V_n[\mathbf{A}^n]. \end{aligned} \quad (84)$$

Therefore, the relativistic interaction Hamiltonian contains single-, double-, and multi-quantum contact vertices. There is no reason to artificially retain only the lowest degrees of the field, as was done in the former QFT. The number of significant terms in the series (84) is determined only by the required order in the interaction constant, the conservation laws, and the kinematics of the process. In particular, when interacting with a spinless particle in the s-wave state, even degrees remain, while in the p-wave state, odd degrees may also exist.

Thus, in TS QFT, the multiparticle contact vertices for bosons follow directly from the expansion of the relativistic Hamiltonian and there are many such ones.

For fermions, the linearization of the relativistic HJ equation leads to the matrix Dirac Hamiltonian:

$$H_D = \boldsymbol{\alpha} \cdot (\mathbf{P} - q\mathbf{A}) + \beta m + qA_0, \quad (85)$$

from which linear fermion vertices follow. At squaring the first-order operator, a spin interaction arises with $\sigma^{\mu\nu} F_{\mu\nu}$.

Thus, the differences in the interactions of bosons and fermions follow from the different realization of the same relativistic mass condition.

5.2. Time-symmetric S-matrix and perturbation theory

To obtain the transition amplitudes, we will use the interaction representation. In this representation, the time dependence of the operator amplitudes determines the free Hamiltonian H_0 , and the evolution of the system's state is determined by the interaction Hamiltonian H_I . If $|\Phi(t)\rangle_I$ is the state in this representation, then we have:

$$i \frac{\partial}{\partial t} |\Phi(t)\rangle_I = H_I(t) |\Phi(t)\rangle_I. \quad (86)$$

For positive energy states, evolution proceeds forward in the ordinary time t , while for negative energy states, it proceeds backward. Therefore, the ordinary chronological ordering operator T_+ is supplemented by the antichronological ordering operator T_- introduced in Section 4. Together, they form a time-symmetric ordering operator $\hat{T} = T_+ T_-$, which acts selectively on the operator amplitudes of each energy sector. The ordinary operator T_+ orders positive energy operators in order of increasing time from right to left, while T_- orders negative energy operators in reverse order.

The evolution operator is therefore written as

$$U(t_2, t_1) = \hat{T} \exp \left[-i \int_{t_1}^{t_2} dt H_I(t) \right]. \quad (87)$$

At passing to free states in the asymptote, the evolution operator yields S-matrix $S = U(+\infty, -\infty)$, or in local representation

$$S = \hat{T} \exp \left[-i \int d^4x H_I(x) \right] \quad (88)$$

and the amplitude of the process $i \rightarrow f$ is $S_{fi} = \langle f | S | i \rangle$. The expansion (88) yields a time-symmetric perturbation theory series:

$$S = 1 + \sum_{n=1}^{\infty} \frac{(-i)^n}{n!} \int d^4x_1 \cdots d^4x_n \hat{T}[H_I(x_1) \cdots H_I(x_n)]. \quad (89)$$

At each order, matrix elements arise between the physical initial and final states of the multiparticle system. The norms of these states are positive, regardless of the sign of energy of the initial states. Negative-energy states evolve backward in time and, after the ZSF interpretation, are described as corresponding positive-energy antiparticles moving forward in time.

As shown in Section 4, the time-symmetric ordering leads to the causal propagators $D_c(p)$ and $S_c(p)$ describing both energy sections.

At presenting the results in terms of the ordinary time, both oppositely directed evolutions are combined into a standard covariant form. This explains why the diagram technique of TS QFT, in its most compact form, often has the same appearance as the former covariant diagram technique.

However, the coincidence of the forms of the final expressions does not imply the coincidence of their physical interpretation. In TS QFT, there is no physical vacuum of quantum fields continuously fluctuating at every point in space. The ground state of a multiparticle system implies the absence of particles, and virtual processes appear only as intermediate states during particle interactions.

Hermiticity of the complete Hamiltonian $H^\dagger = H$ ensures the unitarity of the evolution operator $S^\dagger S = 1$. The usual unitarity relations and the optical theorem follow from this. The intermediate states arising from cutting the diagrams are states of physically possible particles with positive norms, and the sign of energy determines the direction of their evolution.

Thus, the perturbation theory of TS QFT retains the main technical advantages of covariant QFT, but its initial objects are multiparticle states and operator transition amplitudes, rather than the physical vacuum of quantum fields.

5.3. Diagram technique

Decomposing time-ordered products (89) into normally ordered products and then into propagators leads to a diagram technique. The inner lines correspond to the causal propagators of particles, the outer lines describe the actual initial and final states, and the vertices are the transitions between multiparticle states, determined by the terms H_I .

At the vertices, due to translational symmetry, the total 4-momentum is conserved. For internal lines, 4-momentum generally does not satisfy the mass-shell condition, and integration is performed over it. Such lines describe virtual intermediate states.

For bosons, a distinctive feature of the HJ formalism is the emergence of new contact interactions. The expansion (84) shows that, in addition to the usual processes constructed from vertices, contact vertices $V_1, V_2, V_3, V_4, \dots$ also can arise at the same order of charge.

Therefore, the complete matrix element of a given order must contain all diagrams following from the corresponding terms of the same Hamiltonian, including contact contributions following from the term of the Hamiltonian quadratic in the gauge field. It is the sum of the exchange and contact diagrams that forms the physical amplitude and must satisfy the gauge identities. Therefore, contact vertices are not simply additions to restore gauge invariance; in the HJ approach, they arise from the same relativistic Hamiltonian as linear vertices. This is important for the interaction of photons with

charged bosons and for massive vector bosons. As the number of external gauge quanta increases, the exact expression (83) generates vertices of increasingly higher orders. Their role is determined by the order in the coupling constant and the kinematic regime of the process.

The situation is different for a fermionic field. The linearized Dirac Hamiltonian contains the interaction linearly, and therefore multiphoton processes in lower orders are constructed primarily by sequentially connecting such vertices with internal fermionic propagators. The difference between the bosonic and fermionic cases follows from the difference in their relativistic Hamiltonians.

Closed internal lines form loop diagrams. For a loop with internal 4-momentum k , an integral of the form

$$\int \frac{d^4k}{(2\pi)^4} N(k, p_i) \prod_j D_j(k, p_i), \quad (90)$$

arises, where N define the vertices and numerators of the propagators, and D_j are the denominators of the internal lines. The physical meaning of such diagrams in TS QFT differs from the picture of "vacuum fluctuations." A connected loop diagram is part of the amplitude of a specific process and describes a virtual sequence of intermediate multiparticle states between the interaction vertices.

Uncoupled vacuum diagrams do not correspond to transitions between physical initial and final states. Normalized matrix elements do not contain them and, accordingly, they do not contribute to the observed amplitudes. Thus, the diagram technique for interacting particles is consistent with the absence of a ground-state zero-point energy for systems of free particles.

Loop corrections to physical processes are preserved. Physical vacuum fluctuations were consequences of the hypothesis of a physical quantum field in the vacuum, while virtual intermediate states follow directly from quantum particle fluctuations calculated using perturbation theory for an interacting multiparticle system.

The time-symmetric structure also manifests itself within the loop. Integration over k^0 of the causal propagators includes its poles of both energy signs. These contributions correspond to two directions of evolution of intermediate states, combined in a single covariant expression.

At the tree level, this diagram technique leads to the usual rules for constructing amplitudes: external states correspond to wave functions or polarization vectors, internal lines to causal propagators, vertices to matrix elements of the corresponding terms H_I , and at each vertex, energy, momentum, and internal quantum numbers are conserved. At the loop level, these same rules lead to integration over independent internal 4-momenta.

The HJ approach leads to the emergence of contact interactions with the number of fields increasing with each order of perturbation theory. In loop diagrams, this leads to a faster growth of the number of propagators compared to the number of vertices. Therefore, such diagrams themselves are ultraviolet-finite. The observed effects of these interaction terms are discussed in the book [13].

At sufficiently large virtual energies, however, a new circumstance arises that was not taken into account in the former QFT: the large physical energy of the intermediate state, included in energy integral, is the source of the external gravitational field that must be taken into account.

This means that the former covariant diagram technique was, in fact, valid without modification only up to scales where the gravitational field of the intermediate state could be neglected. These are energy regions considerably smaller than the Planck energy. As the Planck scale is approached, the external gravitational field created by the

system, i.e., the curvature of spacetime, invalidates the basic condition for the diagram technique - a flat spacetime background - and one must now follow general relativity.

Thus, TS QFT approach leads to the conventional causal diagram technique as a compact form for describing transitions in a multiparticle system, but it clarifies its physical content. In it, the vertices follow from the full relativistic interaction Hamiltonian, where contact and exchange diagrams of the same order must be considered jointly, and at high energies in the loops, the external gravitational field of the system must be taken into account.

The last point leads to the gravitational self-regularization of loop diagrams, discussed in the next section.

5.4. Gravitational self-regularization and finiteness of TS QFT

In the propagator approach of RQM, the external gravitational field of the particle system in loop diagrams was taken into account at the Planck length scale and the required gravitational dilation of proper times in this field according to the general relativity. This length is the gravitational radius of the system's of Planck energy, Λ_g , and it was shown that gravitational freezing of all processes at this scale leads to gravitational self-regularization and the finiteness of perturbation theory [8].

In TS QFT, perturbation series and diagram techniques are obtained through more sophisticated technics, but they remain essentially the same as in TS RQM propagator approach. Therefore, gravitational self-regularization in TS QFT occurs in the same way as in TS RQM.

Gravitational self-regularization makes the contributions to the S-matrix of all loop diagrams of the perturbation theory series finite. Moreover, interactions with SM gauge fields and quantum gravity are also renormalizable due to smallness of the effective interaction constants at pre-Planck energies.

The self-interaction of a scalar field associated with a scalar particle formerly yielded excessively large (power-law) contributions at Planck scales. In TS QFT, the scalar field remains an external field, and the interactions between scalar particles are not necessarily determined by the self-interaction terms of the external field and require clarification by particle physics. Therefore, in TS QFT, the problems of the former QFT with a scalar field largely disappear and become problems of particle physics.

The finiteness of perturbation theory in TS QFT leads to various and non-trivial consequences:

1. TS QFT is a finite and consistent physical theory.
2. There is complete agreement between TS QFT and general relativity.
3. The long stage of the formation of QFT, when it contained internal contradictions in its foundations, is coming to an end.
4. There is no longer a serious need to search for and test very radical hypotheses, and this fact significantly simplifies the situation in particle physics.
5. Gravitational freezing rules out hypotheses about a "hot" vacuum on small scales, and the structure of the ground state remains simple.
6. Quantum gravity becomes consistent and simpler than expected, since it reduces to contributions at $\Lambda < \Lambda_g$, where it suffices to take one-loop diagrams with gravitons, while the contributions from higher-order diagrams turn out to be very small.
7. SM and quantum gravity can be unified into a general theory, the Extended SM (ESM), with symmetry group $G \times SU(3) \times SU(2) \times U(1)$, where G is the local symmetry group of the graviton field, which is a physically significant part of the group of general coordinate transformations.
8. After the first successful unification in electroweak theory, attempts to further unify SM fields were natural and will remain a major research area in particle physics. Former attempts to unify the three gauge fields into one were based on the hypothesis

that at energies of the order of 10^{15} GeV the charges of the three gauge fields would be equal. However, under gravitational regularization, the curves in the asymptotic behavior of the three charges intersect only pairwise and at three different points, energy scales of which differ by several orders of magnitude. Therefore, the Grand Unification of interactions remains problematic, and new evidence is needed.

9. The situation in particle physics is simplified by the fact that the local TS QFT in 4-dimensional space-time turned out to be sufficient to solve former problems and there is no need for unusual hypotheses about either non-locality (except for Planck scales) or an increase in the number of spatial dimensions.

10. Field models with nontrivial topologies form a special class of models. Some of these, when formulated physically correctly, will be further developed, as accounting for nontrivial nonlinear effects improves the accuracy of calculations and often leads to the discovery of new phenomena.

11. Gravitational regularization opens new prospects for composite models, since, firstly, it predicts the composite nature of the scalar boson, possibly at distances that will be experimentally accessible in the near future. In this regard, the theoretical problems of composite models become particularly pressing, and the main ones will likely be resolved in the near future.

12. Composite particle models will be a promising area for further development in particle physics, also in terms of unifying external interaction fields. Successful composite models would reduce the number of primary particles, which would simplify unification models and limit their versions.

Thus, within the framework of TS QFT, a closer attention to composite models and a concentration of efforts on solving their problems is expected, which, if successful, will lead to a situation, natural from the point of view of historical tradition, when the progress of particle physics will again be associated with a reduction in the number of primary particles, simplification of their models and unification of their interactions.

6. Conclusion

Thus, TS QFT, based on TSR [7] and developing TS RQM [9], is a consistent theory describing in covariant form the quantum properties of relativistic fields. The field operators are not given the meaning of physical field operators; they remain what they actually were in the former QFT – the operator amplitudes of transition probabilities in a multiparticle system. Therefore, they are complex, regardless of the presence of charge. Doubling the state space for two energy signs leads to the fact that relativistic fields in TS QFT describe systems of particles and antiparticles.

For multiparticle systems, the relativistic canonical formalism is most consistent in the HJ formalism, where energy enters the HJ equation to the same degree as 3-momentum. This not only simplifies the theory of free particles but also modifies the theory of bosons interacting with gauge fields. Contact interactions appear, with the number of fields increasing with each order of perturbation theory. Since this leads to a more rapid increase in the number of propagators in loop diagrams compared to the number of vertices, such diagrams are ultraviolet-finite without regularizations. The observable effects of these interaction terms are discussed in the book [13].

Time integration in action functions and perturbation theory occurs in both directions of time, depending on the sign of energy. Therefore, in TS QFT, both forward-time T_+ and backward-time ordering T_- operators are introduced, which are combined into a time-symmetric ordering operator $\hat{T} = T_+ T_-$, which leads to causal propagators and diagram techniques.

The evolution of two types of states in two time directions leads to a modification of the diagram technique. Writing expressions in ordinary time with ordinary limits makes the formulas more compact and reproduces the results of the standard diagram technique. Thus, TS QFT explains why the diagram technique correctly described the observed effects.

TS QFT does not contain the zero-point energy and is consistent with experiments whose results were attributed to the zero-point energy but are in fact explained by contributions from real source fields. Since the absence of a ground-state energy in TS QFT is independent of the masses, supersymmetry breaking in supersymmetric theories also does not change the ground-state energy.

As in TS RQM, there is gravitational self-regularization of loop diagrams at Planck scales, as required by general relativity, making perturbation theory series finite for any interactions. The SM gauge field sector and quantum gravity are also renormalizable due to smallness of the effective coupling constants at pre-Planck energies.

Thus, TS QFT is a new theory, radically different from the former QFT in its physical aspects, but retaining most part of its mathematical apparatus and explaining the reason for its effectiveness despite all the shortcomings of QFT. TS QFT has no divergent ground-state energies or loop correction contributions, nor does it have problems with negative energies and state norms. Moreover, all this is achieved not by introducing new hypotheses and entities, but on the contrary, by eliminating them from the previous QFT and carefully taking into account everything that was previously unreasonably ignored.

A more detailed description of TS QFT and its consequences is given in the book [13].

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