of the type (77). Using then the inequality (76), we obtain an upper bound on the slope for this class of solution:

\[ \frac{\sqrt{S^2}}{M^2} < \frac{2}{\sqrt{m}}. \] (80)

Similar arguments for the solution (73), (74) give the same estimate as for the solution (70), (71), i.e., the inequality (80). Thus, it follows from the inequalities (78) and (80) that in the baryon string model we consider the slope of the classical Regge trajectories is always less than the meson slope \( \frac{\sqrt{S^2}}{M^2} = \frac{1}{2n} \) here, following [1, 2], we assume that the constant \( a' \) in (3) is equal to the corresponding meson constant.

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**LITERATURE CITED**

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**PERTURBATION METHOD IN THE THEORY OF KINETIC EQUATIONS**

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Including in the definitions of correlation functions singular corrections that describe "self-correlations," it is possible to find a rigorous, almost trivial solution of the complete BBGKY hierarchy for a system of charged particles corresponding to motion of them in a self-consistent field. Study of the averaged small deviations from this motion makes it possible to construct a scheme of successive approximations. In this manner, an expansion is obtained for the single-particle distribution function which is equivalent to a generalization of Grad's moment method to the phase space. In the first order of perturbation theory, an approximate Lenard–Balescu equation that differs from the result of its direct linearization is obtained. The proposed approach makes possible a more consistent approximate treatment of statistical systems.

**1. On an Exact Solution of the Complete BBGKY Hierarchy**

The most complete and rigorous description of nonequilibrium statistical systems is achieved by means of the hierarchy of mutually coupled equations for the sequence of distribution functions known as the Bogolyubov chain (BBGKY hierarchy); it can be obtained by systematic partial integration of the Liouville equation [1]. Truncation of the hierarchy by some or other physical assumptions leads to various closed approximate equations for the single-distribution function; these are known as kinetic equations [1, 2].

The aim of the present paper is to construct a method of successive approximation for the complete hierarchy before truncation; for this, it is necessary to know at least one exact solution of it.

Unfortunately, the complete hierarchy does not admit even the trivial solution for the correlation functions. To see this, it is sufficient to take a glance at, say, the second equation of this hierarchy, written down for the binary correlation function \( g(x_1, x_2, t) \) and containing the triple correlation function \( h(x_1, x_2, x_3, t) \). In this equation, it is not possible to set \( g(1, 2) = h(1, 2, 3) = 0 \), since the equation is inhomogeneous: its right-hand side contains a term that is determined by a product of the single-particle distribution functions \( f(x_1, t) \) and does not contain correlation functions (the following equations of the hierarchy have a similar form). It would thus appear that the Vlasov equation [3, 4], which corresponds to the
neglect of correlations and is obtained from the first equation of the hierarchy with \( g(1, 2) = 0 \), contradicts, strictly speaking, the complete system of equations.

However, if in writing down the hierarchy we go over to new unknowns, including self-correlations in the definitions of the \( s \)-particle distribution functions, the equations for such correlation functions, which then become singular, are homogeneous and admit an exact trivial solution, and this can be taken as the basis of a method of successive approximation. Such a transition can be implemented by defining the modified \( s \)-particle distribution functions in a system of \( N \) particles \( \mathcal{f}_s \) \((s=1, 2, \ldots, N)\) in accordance with Klimontovich’s approach [5] as averaged microscopic densities:

\[
\frac{N!}{(N-s)!} \mathcal{f}_s(1, \ldots, s) = \sum_{\alpha_1, \ldots, \alpha_s} \delta(x_{1s} - x_{\alpha_1}(t)) \ldots \delta(x_{s} - x_{\alpha_s}(t)),
\]

where \( x_i = (x_i, v_i) \) characterize the points in the phase \( \mu \) space of the system (sometimes they will be denoted simply by numbers), and \( x_{\alpha_i}(t) \) determines the law of motion of the particle with the number \( \alpha_i \) \((i=1, 2, \ldots, s)\).

In the sum (1), each of the indices \( \alpha_s \) takes all values from 1 to \( N \), and the bar denotes averaging over all possible initial states \( \Gamma = \{x_{01}, \ldots, x_{0N}\} \) (by virtue of Liouville's theorem, any instant of time can be taken as the "initial" instant). If the system contains particles of different species, the corresponding functions (1) are defined for each species separately.

The distribution functions \( \mathcal{f}_s \) differ from the functions \( f_s \), usually considered in that the sums (1) which define them include terms with all summation indices \( \alpha_s \), including coincident indices, whereas in the definition of \( f_s \) terms with the same \( \alpha_s \) are omitted (which is usually indicated by a prime on the summation sign). Such a modification obviously does not affect the single-particle distribution function \( \mathcal{f}_1 \), and the modified correlation functions \( \mathcal{g}_s, \mathcal{h}_s \), etc., can be expressed in terms of the ordinary ones as follows (for \( N \gg 1 \)):

\[
\frac{\mathcal{f}_s(t)}{2} = \frac{f_s(t)}{N}, \quad \frac{\mathcal{g}(1, 2)}{2} = \frac{g(1, 2)}{N}, \quad \frac{h(1, 2, 3)}{3} = \frac{h(1, 2, 3)}{N} + \frac{3}{N} [g(1, 2)g(2, 3) + g(2, 3)g(3, 1) + g(3, 1)g(1, 2)],
\]

whereas the inverse transformation has the form

\[
\frac{f_s(t)}{2} = \frac{\mathcal{f}_s(t)}{N}, \quad \frac{g(1, 2, 3)}{3} = \frac{h(1, 2, 3)}{N} + \frac{3}{N} [g(1, 2)g(2, 3) + g(2, 3)g(3, 1) + g(3, 1)g(1, 2)].
\]

The singular functions \( \mathcal{f}_s \) have been used for statistical description before now (see, for example, [6]), but the fact that they make it possible to find an exact solution of the complete Bogolyubov chain was not noted. We note also that the exact microscopic solutions of the Vlasov and Boltzmann-Enskog equations found in [4, 7] also have a singular nature (admittedly, for a different reason), but in contrast to (1) the single-particle distribution function is already singular in [4, 7].

Substituting (3) in the first two equations of the hierarchy, we obtain

\[
\left( \frac{\partial}{\partial t} + L_1 \right) f(x, t) = \frac{N}{m} \frac{\partial q(1, 2)}{\partial x_1} \mathcal{g}(1, 2, t) dx_1,
\]

\[
\left( \frac{\partial}{\partial t} + L_1 + L_2 \right) g(x, x, t) = \frac{N}{m} \left[ \frac{\partial f(1, t)}{\partial x_1} \frac{\partial q(1, 3)}{\partial x_1} \mathcal{g}(1, 3, t) + \frac{\partial f(2, t)}{\partial x_2} \frac{\partial q(2, 3)}{\partial x_2} \mathcal{g}(2, 3, t) \right] dx_1 + \frac{N}{m} \left[ \frac{\partial q(1, 3)}{\partial x_1} \frac{\partial f(1, t)}{\partial x_1} + \frac{\partial q(2, 3)}{\partial x_2} \frac{\partial f(2, t)}{\partial x_2} \right] \mathcal{h}(1, 2, 3, t) dx_1.
\]

We have here used the following notation for the linear differential operator with self-consistent field:

\[
L_n = \frac{\partial}{\partial x_n} + \frac{1}{m} \int F^e(r_\xi) - N \int \frac{\partial q(r_\xi - r')}{\partial x_n} f(r', v', t) dr' dv',
\]

where \( F^e \) is the external force acting on the system of \( N \) particles that each have mass \( m \), their interaction being described by the two-body symmetric potential \( q(a, b) = q(|r_a - r_b|) \).

Transforming the equations of the hierarchy, it is necessary to bear in mind that by virtue of the symmetry of the \( \delta \) function,