As a consequence, we obtain from this that functions \( v(x) \) of the form \( u_0(x) + u^2(x) \), and only such functions, are virtual potentials in the Schrödinger equation (2) if \( u(x) \) is an absolutely continuous function satisfying the summability conditions (4) and (5) for \( c_1 = c_2 = 0 \).

Remark. If \( c_1^2 = c_2^2 \), then part 1 of Theorem 1 must be replaced as follows:

1) the function \( r_2(k) \) is continuous for all \( k \in \Gamma \) and satisfies the conditions

\[
r_2(-k) = r_2(k), \quad r_2(k-i0) = r_2(k+i0), \quad -1 \leq r_2(k) < 1, \quad k \in \Gamma, \quad r_2(k) = O(k^{-1}), \quad k \to \pm \infty.
\]

At the same time, the function \( r_2(\pm |c_2|) = -1 \) in the generic situation and \( -1 < r_2(\pm |c_2|) < 1 \) in the virtual case, i.e., when \( W[f(x), f(x)] = 0 \) for \( k = \pm |c_2| \).

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LITERATURE CITED

obtain an equation of Lippmann-Schwinger type equivalent to the Low equation. In this paper, we give a different method for linearizing the field-theoretical Low equation; it enables us to eliminate in the linearized equation the terms with a bound deuteron intermediate state and obtain completeness and orthonormalization conditions on the scattering wave functions. An important difference between the linearized equations for NN and πN scattering is that the Low-type equations for NN scattering linearize completely, while the linearized equations for πN scattering will contain in the potential a nonlinear u-channel term associated with a crossing permutation of the π mesons [7]. In order to compare the linear relativistic equations for NN scattering obtained in this paper with the quasi-potential equations of [8–13], we propose in Sec. 3 quasipotential equations whose solution in the off-energy-shell region is identical to the required solution of the Low-type equations for NN scattering.

1. Connection Between Low-Type Equations and
the Medvedev-Polivanov Equations

In the S-matrix formulation of local field theory [2,3], an important part is played by the radiation operators, which for the nucleon Fermi operators in which we are interested have the form

$$\mathcal{J}(x) = -i \mathcal{P} \frac{\delta \mathcal{J}}{\delta \psi^{\text{in}}(x)}; \quad \mathcal{J}(x) = i \mathcal{P} \frac{\delta \mathcal{J}}{\delta \psi^{\text{in}}(x)},$$

where $\mathcal{J}$ is the operator of the S matrix, and $\psi^{\text{in}}(x)$ is the nucleon operator in the asymptotic in state. In the S-matrix formulation [2,3], the nucleon quantum field $\psi(x)$ is determined by means of an equation of Yang-Feldman type:

$$\psi(x) = Z_2 \psi^{\text{in}}(x) + \int S^{\text{ret}}(x-y) \mathcal{J}(y) dy,$$

where $S^{\text{ret}}(x-y)$ is the retarded Green's function of the spinor field, and $Z_2$ is the renormalization constant of this field [2,14]. Further, we introduce the operator $b_{p_s}(x)$, which in accordance with the asymptotic Lehmann-Symanzik-Zimmermann condition [14] goes over into $b_{p_s}(\in)$ in the corresponding weak limit $x^0 \to -\infty$:

$$b_{p_s}(x) = Z_2^{-1} \int d^4x S(x) \gamma^\mu e^{-ip_s \gamma^\mu} u(p_s),$$

where $p=(1m^2+p^2, p)$ is the nucleon 4-momentum, and $u(p_s)$ is the well-known spinor function. It follows from the definitions (3) and (2) that

$$b_{p_s}(x) = b_{p_s}(\in) - \frac{i}{Z_2} \int \mathcal{J}(y) e^{-ip_s \theta(x-y)} dy u(p_s).$$

From Eq. (4), we can readily obtain the following identity for the equal-time anticommutators of the operators $b_{p_s}(\in)$ and $\mathcal{J}(x)$:

$$\{\mathcal{J}(0), b_{p_s}(\in)\} = \{\mathcal{J}(0), b_{p_s}(\in)\} - \frac{i}{Z_2} \int d^4y \theta(-y^0) e^{-ip_s \theta}(\mathcal{J}(0)) \mathcal{J}(y) u(p_s),$$

this being the basis for the derivation of Low-type equations in field theory. In particular, in our case of nucleon-nucleon scattering, we must average the identity of (5) over the single-nucleon states of the in and out sets, after which we obtain

$$\hat{\mathcal{J}}_{x_x} \langle p_s' s'_s | \mathcal{J}_{p_s' s'_s}(0) | p_s s_s \rangle_{\text{in}} = \hat{\mathcal{J}}_{x_x} \langle p_s' s'_s | \{\mathcal{J}_{p_s' s'_s}(0), b_{p_s}(\in)\} | p_s s_s \rangle_{\text{in}} +$$

$$+ \int d^4y \langle p_s' s'_s | T(\mathcal{J}_{p_s' s'_s}(0)) \mathcal{J}(y) | p_s s_s \rangle_{\text{in}} e^{-ip_s \theta},$$

where $\hat{\mathcal{J}}_{x_x}$ is the operator of antisymmetrization of the identical nucleons, expressed in terms of the nucleon transposition operator $\hat{P}_n$: $\hat{P}_n=\frac{1}{2}(1-P_n)$,

$$\mathcal{J}_{p_s}(x) = Z_2^{-\frac{1}{2}} u(p_s) \mathcal{J}(x).$$

Further, using the integral representation of the step functions, we insert

$$\hat{1} = \sum_n |n; \in \rangle \langle n; \in |$$

in the second term on the right-hand side of Eq. (6) between the operators $\mathcal{J}_{p_s'}$ and $\mathcal{J}_{p_s}$. After this, to obtain a closed system of equations in the complete set...

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