TABLE I. New limits on the configuration mixing in the GDR of ^{12}C obtained with the (p,γ_0) reaction using polarized protons, as compared with the old limits derived from the unpolarized reaction.

Data	δ	$\min\|s_{1/2}\|^2$	$\max\|d_{3/2}\|^2$	$\min\|d_{5/2}\|^2$
Unpolarized	0	0.01	0.52	0.10
Polarized included	$-15°$	0.08	0.14	0.23
	0	0.06	0.26	0.21
	15°	0.02	0.49	0.07

from $\delta=0$, the paths for $d_{3/2}|d_{3/2}|$ and $d_{5/2}|d_{5/2}|$ change substantially for solution II, but there is little change for solution I even for δ as large as $\pm 90°$. The limits on the minimum values for the $s_{1/2}$ and $d_{5/2}$ contributions and the maximum value for the $d_{3/2}$ contribution are shown in Table I and compared with the previous limits based only on the unpolarized data.[2]

If the wave function of the GDR approximates a bound-state wave function, the phases of the T-matrix elements are just the Coulomb phase shifts for the initial channel: $\delta=0$, $\varphi_d-\varphi_s=27°$ at $E_p=6$ MeV and $\varphi_d-\varphi_s=18°$ at $E_p=14$ MeV. It is interesting to note that such values are allowed in solution I. Departure from the Coulomb phase shift reflects the extent to which the GDR wave function departs from the bound state form. We note, however, that a Coulomb plus hard-sphere phase shift, as is usually assumed in the literature,[8] produces $\delta=0$ and $\varphi_d-\varphi_s\simeq 100°$ which is not allowed in either solution I or II.

An analysis of the γ_1 radiation is considerably more complex than that for γ_0, because levels with $J=1^-$, 2^-, and 3^- can each produce $E1$ radiation to the 2^+ first excited state of ^{12}C. If the $d_{3/2}$ and $d_{5/2}$ phases are nearly equal, then it is not possible to explain the observed asymmetry with a single 1^- or 3^- level. However, a single 2^- level or interferences such as $(1^-,3^-)$, $(2^-,3^-)$, and $(1^-,2^-)$ are able to produce the observed asymmetry.

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Covariant Diastrophic Quantum Field Theory

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Elementary arguments determine the operator superstructure for a special but broad class of covariant field theories. Field operators and generators of interest are shown to be bilinear expressions in conventional creation and annihilation operators. Nontrivial interaction is incompatible with canonical (anti) commutation relations.

Choose any covariant quantum field theory that you like; augment the configuration space variables $\vec{x}$ by an auxiliary real variable w; insist on dynamical independence for all space and time of fields with distinct w values: The result is a covariant "diastrophic" quantum field theory, a covariant model having several remarkable properties some of which we report in this Letter.

By way of illustration, consider the diastrophic scalar field theory formally characterized by the Langrangian density

$$\mathcal{L}=\tfrac{1}{2}[\partial_\mu\varphi(x,w)]^2-\tfrac{1}{2}m_0^2\varphi^2(x,w)-V(\varphi(x,w)), \quad (1)$$

where the derivatives ∂_μ act only on the space-time variables x^μ. The formal equation of motion,

$$(\Box+m_0^2)\varphi(x,w)=-V'(\varphi(x,w)), \quad (2)$$

is no less divergent than in the "base" theory (no w), and clearly such models are best approached outside of perturbation theory.[1]

The fundamental symmetry that we exploit is complete independence of the field for distinct values of $w\in R$, a property we call ultralocality.[2] Ultralocality enables the operator "superstructure" of the diastrophic quantum field theory to be completely determined. This superstructure is summarized in theorems I and II below and provides severe restrictions on any such theory. Although (neutral) scalar fields are used for illustration, the arguments extend to higher spins, to spinor fields (briefly discussed), to coupled fields, etc.

How does ultralocality manifest itself? Let $|0\rangle$ denote the vacuum and $\varphi(x,w)$ denote the local field, which we suppose becomes a self-adjoint operator when smeared with a suitable real test function $f(x,w)$. Then ultralocality implies that

$$\langle 0|\exp[i\textstyle\iint\varphi(x,w)f(x,w)\,dx\,dw]|0\rangle =\exp[-\textstyle\int L\{f(\cdot,w)\}\,dw] \quad (3)$$

for some functional L. From this single fact it follows that φ is the analog of an infinitely divisible random variable,[3,4] and that φ may be realized as a bilinear expression in conventional (Fock representation) creation and annihilation operators.[5,6] In this note we demonstrate this operator realization directly by making use of the simple, but powerful, fact that smearing the product of two (or more) formal creation operators taken at the same point does not lead to an operator.

Let $\{A_l(w): w\in R, l\in Z\}$ denote a conventional set of commuting annihilation operators, and let $|0\rangle$ be the unique state for which $A_l(w)|0\rangle=0$ for all w and l. The formal creation and annihilation operators fulfill

$$[A_l(w),A_{l'}^\dagger(w')]=\delta_{ll'}\,\delta(w-w'), \quad (4)$$

and we denote by $\mathfrak{H}$ the usual Hilbert space spanned by repeated action of the $A^\dagger$ on $|0\rangle$. Then, for any diastrophic scalar theory we have the following theorem.

Theorem I: There exists a c-number Hermitian matrix field $\Lambda_{ll'}(x)$, a c-number vector field $\overline{\Lambda}_l(x)$, and a real number C such that the local diastrophic quantum field operator is given by

$$\varphi(x,w)=\sum_{l,l'}A_l^\dagger(w)\Lambda_{ll'}(x)A_{l'}(w)+\sum_l[A_l^\dagger(w)\overline{\Lambda}_l(x)+\overline{\Lambda}_l^*(x)A_l(w)]+C. \quad (5)$$

Proof of theorem I.—In the Hilbert space $\mathfrak{H}$ all operators, including the field, are given as functions of the basic set $A_l(w), A_l^\dagger(w)$ for all $w\in R$, $l\in Z$. Ultralocality forces the local field $\varphi(x,w)$ to be constructed from elements of the basic set *at the point* w. The most general expression of this type leading to a local, formally self-adjoint operator is bilinear in $A^\dagger$ and A as in (5).

The operator realization (5) permits a ready calculation of the truncated vacuum expectation values (TVEV): Specifically, $\langle 0|\varphi(x,w)|0\rangle^T=C$; while for $n\geq 2$,

$$\langle 0|\varphi(x_1,w_1)\varphi(x_2,w_2)\cdots\varphi(x_n,w_n)|0\rangle^T=\delta(w_1-w_2)\delta(w_2-w_3)\cdots\delta(w_{n-1}-w_n)$$
$$\times\sum_{l_1,\cdots,l_{n-1}}\overline{\Lambda}_{l_1}^*(x_1)\Lambda_{l_1l_2}(x_2)\cdots\Lambda_{l_{n-2}l_{n-1}}(x_{n-1})\overline{\Lambda}_{l_{n-1}}(x_n). \quad (6)$$

Evidently the functional L in Eq. (3) is implicitly given by the TVEV.

Several general properties are immediate from (5) and (6): Physically interesting fields have $\overline{\Lambda}_l(x)\not\equiv 0$; quasifree fields have $\Lambda_{ll'}(x)\equiv 0$; fields with nonvanishing higher-order TVEV require $\Lambda_{ll'}(x)\not\equiv 0$; fields fulfilling canonical commutation relations (CCR),

$$[\varphi(\vec{x},w),\dot{\varphi}(\vec{x}',w')]=i\delta(\vec{x}-\vec{x}')\delta(w-w') \quad (7)$$

at any equal time, require $\Lambda_{ll'}(x)\equiv 0$, and hence are quasifree.

From (5) it follows that any asymptotic field $\varphi_{\rm ex}(x,w)$ must also have a bilinear structure, i.e.,

$$\varphi_{\rm ex}(x,w)=\sum_{l,l'}A_l^\dagger(w)\Lambda_{ll'}{}^{\rm ex}(x)A_{l'}(w)+\sum_l[A_l^\dagger(w)\overline{\Lambda}_l{}^{\rm ex}(x)+\overline{\Lambda}_l{}^{*\,\rm ex}(x)A_l(w)]. \tag{8}$$

If $\varphi_{\rm ex}(x,w)|0\rangle$ yields stable one-particle states, then $\overline{\Lambda}_l{}^{\rm in}(x)=\overline{\Lambda}_l{}^{\rm out}(x)\equiv\overline{\Lambda}_l{}^0(x)$. Scattering arises only when $\Lambda_{ll'}{}^{\rm ex}(x)\not\equiv 0$, and $\Lambda_{ll'}{}^{\rm in}(x)\not\equiv\Lambda_{ll'}{}^{\rm out}(x)$; hence imposing CCR leads to no scattering. The S matrix has the form

$$S=\exp[i\sum_{l,l'}\int A_l^\dagger(w)\eta_{ll'}A_{l'}(w)\,dw], \tag{9}$$

where $\eta_{ll'}$ is an appropriate Hermitian matrix annihilating $\overline{\Lambda}_l{}^0$ and connecting $\Lambda_{ll'}{}^{\rm in}$ and $\Lambda_{ll'}{}^{\rm out}$ in an obvious fashion.

Besides the field, *all* fundamental generators of interest are bilinear expressions as well. Let $\mathcal{G}$ denote a generic generator of the diastrophic theory (i.e., Hamiltonian $\mathcal{H}$, space-translation generator $\mathcal{P}$, boost generator $\mathcal{K}$, dilatation generator for scale-invariant theories, etc.). Then we may assert a second theorem.

Theorem II: For each fundamental self-adjoint generator $\mathcal{G}$ of a diastrophic theory, with $\mathcal{G}|0\rangle=0$, there exists an associated Hermitian matrix $\mathcal{G}_{ll'}$ such that

$$\mathcal{G}=\sum_{l,l'}\int A_l^\dagger(w)g_{ll'}A_{l'}(w)\,dw. \tag{10}$$

Proof of theorem II.—As a construct of the basic operator set that respects ultralocality, each fundamental generator $\mathcal{G}$ necessarily has a bilinear form (essentially theorem I). Self-adjointness plus the condition $\mathcal{G}|0\rangle=0$ eliminates both constant and linear constituents, leaving (10) as the only possibility.

If $\mathcal{G}$ denotes a generator of the diastrophic theory, it is plausible to identify $g_{ll'}$ as the corresponding generator of the base theory. For example, if $\mathcal{G}^k$, $k=1,\cdots,K$, is a set of generators that form a Lie algebra (e.g., Poincaré), the matrices $g_{ll'}{}^k$ satisfy the same Lie algebra. However, there are unexpected differences in spectral characteristics. For example, the Hamiltonian

$$\mathcal{H}=\sum_{l,l'}\int A_l^\dagger(w)h_{ll'}A_{l'}(w)\,dw \tag{11}$$

satisfies $\mathcal{H}|0\rangle=0$, and will be nonnegative, $\mathcal{H}\geq 0$, provided $\{h_{ll'}\}\equiv h\geq 0$. But a nondegenerate ground state requires $h>0$ so that the base field theory has no ground state.

In order for the (bilinear!) Hamiltonian (e.g., the interaction potential) to involve various powers of the (bilinear!) field operator φ, infinite renormalizations must enter in. These renormalizations have been worked out in detail for simpler but related cases,[2] and the techniques in the present case are similar.[6] For present purposes we merely note that infinitely renormalized products of the local field $\varphi(x,w)$—all factors taken at the same w—do exist that *still* have a bilinear expression in the basic creation and annihilation operators.

For fermions only a few changes are involved. Consider a local, diastrophic spinor field operator $\psi(x,w)$. We take independence of distinct w values to imply anticommutativity of such fields, namely,

$$\{\psi(x,w),\psi(x',w')\}=0, \tag{12}$$

whenever $w\neq w'$. Besides the basic operators $A_l(w)$ and $A_l^\dagger(w)$ introduced above we consider an additional, anticommuting set $a_l(w)$ and $a_l^\dagger(w)$ fulfilling $a_l(w)|0\rangle=0$ and

$$\{a_l(w),a_{l'}^\dagger(w')\}=\delta_{ll'}\delta(w-w'),\quad \{a_l(w),a_{l'}(w')\}=\{a_l^\dagger(w),a_{l'}^\dagger(w')\}=0. \tag{13}$$

We now define $\mathfrak{H}$ as the space spanned by repeated action of both $A^\dagger$ and $a^\dagger$ on $|0\rangle$. Ultralocality restricts $\psi(x,w)$ to come from elements of the augmented basic set at the point w. The most general expression (with both ψ and $\psi^\dagger$ local operators) is bilinear; however, (12) restricts this to the form

$$\psi(x,w)=\sum_{l,l'}[A_l^\dagger(w)M_{ll'}(x)a_{l'}(w)+a_{l'}^\dagger(w)N_{ll'}{}^*(x)A_l(w)]+\sum_l[a_l^\dagger(w)\overline{M}_l(x)+\overline{N}_l{}^*(x)a_l(w)] \tag{14}$$

for suitable c-number spinor field coefficients. Canonical anticommutation relations for ψ require $M_{ll'}(x)=N_{ll'}(x)\equiv 0$, which is just the condition leading to vanishing higher-order TVEV. Scattering proceeds as in the scalar case, and fundamental generators are again bilinear, having the generic form

$$\mathcal{G} = \sum_{l,l'} \int [a_l^\dagger(w) \bar{g}_{ll'} a_{l'}(w) + A_l^\dagger(w) \hat{g}_{ll'} A_{l'}(w)]\, dw \qquad (15)$$

for suitable Hermitian coefficients.

We have seen that the power of ultralocality is far reaching. Nowhere did our arguments make use of the covariance of the base theory, nor of the dimensionality of space. Thus nearly every theory has similar properties. The original ultralocal theories are just single degree of freedom base theories, and are by now rather well under control[2,6]; for example, even for such models, interacting theories cannot have CCR. Hopefully, the study of such models will assist the study of genuine covariant diastrophic fields.

The solution of a covariant diastrophic field such as outlined in this Letter would be interesting on at least two counts. First, it would be a true covariant theory having infinite mass, field-strength, and coupling-constant renormalizations[6]; but, second, and more important, such theories may relate very closely to their base theories. For example, the *classical* solution to an equation like (2) is just $\varphi_{\rm cl}(x,w) = \varphi_{\rm cl}{}^w(x)$, namely a w-parametrized set of solutions of the base theory. In addition, one should not overlook the fact that every conventional covariant theory becomes a covariant diastrophic theory (in one less space dimension) if just *one* of the spacial gradient terms is dropped from the Lagrangrian. With this direct connection in mind it would seem not unreasonable if interacting, covariant field operators bore a closer resemblance to the bilinear diastrophic form rather than to the manifestly inequivalent linear form of a quasifree theory.

[1]Diastrophic quantum field theories were introduced in the author's 1971 Boulder lectures [J. R. Klauder, "Functional Techniques and Their Application in Quantum Field Theory," lectures given at the Fourteenth Annual Summer Institute for Theoretical Physics, University of Colorado, Boulder, Colorado, June 21–August 13, 1971 (unpublished)]. The terminology is meant to reflect the extension of a given (base) theory in a new direction (loosely, a diastrophism) parametrized by w.

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Interstellar CN Excitation at 2.64 mm

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A sensitive search was made for 2.64-mm line emission from a cloud of CN, whose excitation is known from optical measurements, with essentially a null result. This provides strong support for the proposition that the excitation temperature deduced from the optical CN lines is equal to the temperature of the microwave background.

The first confirmation of the discovery of microwave background radiation[1] was based upon the optically observed excitation of interstellar CN.[2,3] It had long[4] been known from optical-absorption line ratios that a considerable fraction of interstellar CN radicals is found in the first