Observations on the $T \ln R$ term in the quark-antiquark free energy

Joe Kiskis
Department of Physics, University of California at Davis, Davis, California 95616
(Received 25 November 1985)

Consider the response of a pure gauge theory at temperature T to an external quark-antiquark pair separated by R. In the confining phase, the leading term in the free energy at large R is σR. A string-model calculation has given $T \ln R$ for the next-to-leading term. In this paper, the origin of the $T \ln R$ term is considered in a more general context that includes the analog spin model and the lattice gauge theory at strong coupling. The connection with transverse fluctuations is emphasized.

I. INTRODUCTION

This paper is concerned with pure, non-Abelian gauge fields at finite temperature. It is instructive to consider the response of the system to external disturbances in the form of heavy-quark and -antiquark sources. These are represented by Wilson line operators at fixed positions. It is assumed that the theory confines quarks for a range of temperatures above zero. Throughout this paper, we work in the confining phase.

The line-line correlation function manifests confinement by decaying to zero exponentially fast. If R is the distance between the lines, β the inverse temperature, and σ the finite-temperature string tension, then a finite-temperature version of the area law is

$$C(R) \sim e^{-\sigma\beta R} . \tag{1.1}$$

The corresponding contribution to the quark-antiquark free energy F is σR.

A string-model calculation[1] has shown that there is a power-law correction to (1.1) so that, at large R,

$$C \propto \frac{1}{R} e^{-\sigma\beta R} \tag{1.2}$$

and

$$F = \sigma R + T \ln R . \tag{1.3}$$

This paper will show that the same correction arises in the analog spin model for the lines and in the lattice gauge theory at strong coupling. In each case, it is associated with transverse fluctuations.

The $T \ln R$ term in the free energy is universal in the sense that it does not depend upon the details of the theory. In particular, it is insensitive to the gauge group and the coupling as long as one remains in the confining phase. In this way, it is similar to the large-distance R^{-1} term in the zero-temperature, quark-antiquark potential.[2]

As a consequence, the $T \ln R$ term does not appear in straightforward weak- or strong-coupling expansions of $C(R,g^2)$. The weak-coupling expansion is not valid at large R. Lattice strong coupling is better in that it corresponds to large physical separations. However, the roughening effects that will turn out to be important are not incorporated in the usual strong-coupling expansion. A more general approach is indicated.

The sections that follow present three different views. Section II recalls the relationship of the line-line correlation function to the spin-spin correlation function of a spin system in the same spatial dimension. This gives the first indication that a power correction may be present in the gauge theory. Section III comments on the string-model calculation. Section IV looks at the strongly coupled lattice theory. In this context, the power correction follows from the spectral representation of C plus either approximate rotational invariance or roughness. The rotational-invariance argument admits the possibility that the power of R is less than minus one and that the coefficient of $T \ln R$ is greater than one. However, the other more specific approaches indicate that this is not likely to be the case.

Each argument reveals in a slightly different way that the power correction is a consequence of the transverse fluctuations of the flux tube. The result depends upon the number of transverse spatial dimensions into which it can move. The power R^{-1} is correct for our world of three spatial dimensions. In d spatial dimensions the power is $-(d-1)/2$. The Ornstein-Zernike decay[3] of correlations in spin systems has the same power and analogous physics.

II. SPIN MODELS

Evidence for the power-law correction can be found in a consideration of the relationship of the finite-temperature gauge theory to a spin model. While these indications are suggestive, they are far from conclusive. Stronger arguments are given in Secs. III and IV.

It has been pointed out that Wilson lines in finite-temperature lattice gauge theories are in some ways similar to the spins of lattice spin models.[4] More specifically, the line-line correlation function of SU(N) lattice gauge theory is related to the spin-spin correlation function of a Z_N spin system of the same spatial dimension. For example, SU(2) gauge theory in 3 + 1 dimensions is related to the Ising model in three dimensions. The confining phase of the gauge theory corresponds to the high-temperature phase of the Ising model.

The on-axis correlation function of the Ising model has an Ornstein-Zernike decay:[5]

$$\frac{1}{I}e^{-MI} . \tag{2.1}$$

Since this is a general property of the model that does not depend upon the details of the interactions, we might expect that the simple area law

$$e^{-\sigma a^2 N_T I} \tag{2.2}$$

for the line-line correlation function should be corrected to

$$C(I) \propto \frac{1}{I} e^{-\sigma a^2 N_T I} . \tag{2.3}$$

The continuum limit of this gives (1.2) and (1.3).

One may inquire as to the origin of the I^{-1} factor in the spin model. For our purposes, it will be most useful to note that it has been explained in terms of the transverse fluctuations of random walks that connect the two sites in the correlation function.[6]

III. STRING MODEL

The string picture[7] is particularly well suited to the problem. It automatically selects the confining phase, and the fixed-coupling, large-R limit is straightforward. This section reviews the string-model calculation of the quark-antiquark free energy in which the next-to-leading $T \ln R$ term appears.

Consider a string in three spatial dimensions with ends fixed at (0,0,0) and $(0,0,R)$. Finite-temperature effects are incorporated by applying periodic boundary conditions in a fourth dimension of length β. The string is described by giving its two-dimensional transverse displacement $\mathbf{x}_\perp$ as a function of z and τ:

$$\begin{aligned} &\mathbf{x}_\perp(\tau+\beta,z)=\mathbf{x}_\perp(\tau,z) , \\ &\mathbf{x}_\perp(\tau,0)=\mathbf{x}_\perp(\tau,R)=0 . \end{aligned} \tag{3.1}$$

The action is taken from Ref. 8:

$$S=\sigma\beta R+\frac{\sigma}{2}\int d\tau \int dz(\dot{\mathbf{x}}_\perp\cdot\dot{\mathbf{x}}_\perp+\mathbf{x}'_\perp\cdot\mathbf{x}'_\perp) . \tag{3.2}$$

The free energy is obtained from the logarithm of the functional integral

$$\begin{aligned} C(R) &\propto \int Dx_\perp e^{-S} \\ &\propto e^{-\sigma\beta R}[\det(-\partial_\tau{}^2-\partial_z{}^2)]^{-1} . \end{aligned} \tag{3.3}$$

The determinant with these boundary conditions has been calculated.[1] It gives the leading terms

$$\begin{aligned} C(R) &= \exp\left[-\sigma\beta R+\frac{\pi}{3}\frac{R}{\beta}-\ln\frac{R}{\beta}\right] , \\ F &= \left[\sigma-\frac{\pi}{3}T^2\right]R+T\ln R . \end{aligned} \tag{3.4}$$

We have evaluated the determinant by a different method and obtained the same result. We find that the $T \ln R$ term comes entirely from the modes of the surface that are independent of τ. Since all fixed-τ slices of these surface modes are the same, they are equivalent to the transverse fluctuations of a one-dimensional object in a three-dimensional space.

IV. LATTICE MODEL AT STRONG COUPLING

As noted already, a straightforward strong-coupling expansion of the line-line correlation function does not include the power correction. However, two general properties of the theory that hold at strong coupling together imply that the power correction is present.

The first property is that the lattice line-line correlation function $C(\mathbf{I})$ has a spectral representation with a positive weight. Except for the lack of rotational invariance, this is the same as the textbook[9] result for continuum theories.

The second ingredient relates to the rotational properties of $C(\mathbf{I})$. Let us restrict our attention to the behavior of C on the 3-axis $C(0,0,I_3)$. The power-law correction to the exponential decay of this quantity at large I_3 is closely related to the behavior of $C(\mathbf{I}_\perp,I_3)$ as $\mathbf{I}_\perp$ increases. Both are determined by the spectral density of the 3-axis transfer matrix at threshold.

Two approaches are discussed. The first simply assumes that $C(\mathbf{I})$ is approximately rotationally invariant at large I. It applies to the region of the (N_T,g^2) plane where this is the case. The second uses roughness, which is a crude sort of nearly on-axis rotational invariance, that is present at any finite temperature.

A. Spectral representation

Consider a lattice gauge theory with N_S sites in each of three spatial directions and N_T sites in a fourth imaginary time or inverse temperature direction. Take periodic boundary conditions in all directions and $N_S \to \infty$. Keep N_T finite for a dimensionless temperature $aT=N_T{}^{-1}$. The bare coupling is g^2. If $\mathbf{I}$ labels positions on the spatial lattice and $L(\mathbf{I})$ is a Wilson line at position $\mathbf{I}$, the correlation function is

$$C(\mathbf{I})=\langle L^\dagger(\mathbf{I})L(0)\rangle . \tag{4.1}$$

We seek an expression for the on-axis behavior of C at large separation. The 3-axis is arbitrarily selected.

Begin the spectral decomposition of C by introducing the 3-axis transfer matrix. The discussion that follows is quite familiar if one imagines interchanging the names of the 3 and 4 axes. The result is a zero-temperature Euclidean theory with finite extent and periodic boundary condition in one spatial direction.

The steps[10] in obtaining a spectral representation for C are standard: Work in the gauge where the group elements on links in the 3-direction are unity. Show that the functional integral with the Wilson action is a product of transfer matrices in the 3-direction. Introduce a Hilbert space of configurations at fixed I_3 and, in that, a basis that diagonalizes the 3-axis transfer-matrix and the translation operators in the 1-, 2-, and 4-directions.

Let $\{|\psi\rangle\}$ be this basis. The 1,2-translation eigenvalues of $|\psi\rangle$ are $e^{i\mathbf{P}_\psi\cdot\mathbf{I}_\perp}$ and the transfer-matrix eigenvalues are λ_ψ. $|\,\text{vac}\rangle$ is the translationally invariant state with the largest eigenvalue λ_0. For the other eigenvalues, we write

$$\lambda_\psi=\lambda_0 e^{-E_\psi} . \tag{4.2}$$

Since the Wilson line and all other operators that appear are invariant under 4-direction translations, states with nontrivial behavior under these translations make no contribution. That leaves sums and integrals over positions and momenta in the 1,2-transverse plane.

After the usual manipulations, we are left with the representation

$$C(\mathbf{I}_\perp,I_3)=\int d^2P_\perp \int dE\, e^{iP_\perp\cdot I_\perp}e^{-EI_3}\rho(\mathbf{P}_\perp,E) . \tag{4.3}$$

It is crucial that the weight

$$\rho=\sum_\psi \delta(E-E_\psi)\delta^2(P_\perp-P_\psi)\,|\langle\psi|L|\,\text{vac}\rangle|^2 \tag{4.4}$$

is, for each $P_\perp$, a positive measure in E that vanishes for $E<M$. There is such a positive M because we are in the confining phase. Its value and dependence upon g^2 play no role in this work. The representation (4.3) also reveals the important fact that the large-I_3 behavior of C is determined by the functional form of the spectral weight at threshold.

B. Rotationally invariant forms

This discussion is valid for the region of the (N_T,g^2) plane that is confining and approximately rotationally invariant. For this case, the argument is simple and standard.

Begin by assuming that C has the approximately invariant form

$$C=\frac{1}{I^\alpha}e^{-MI} \tag{4.5}$$

at large $I\equiv|\mathbf{I}|$. M is small and fixed. This large-I behavior of C is fixed by the strength of the singularity in its Fourier transform at $\mathbf{P}^2=-M^2$. Our strategy is to compute the Fourier transform of (4.5), express it in terms of a spectral weight, and compare with (4.3) and (4.4). Consistency will require that α be greater than or equal to one.

The singularity of the Fourier transform at $\mathbf{P}^2=-M^2$ is adequately approximated by replacing the sum on $\mathbf{I}$ by an integral

$$\tilde{C}(\mathbf{P})=\int d^3x\, e^{iP\cdot x}C(x) . \tag{4.6}$$

If this takes the form

$$\tilde{C}=\int d\sigma^2\frac{R(\sigma^2)}{P^2+\sigma^2} , \tag{4.7}$$

then an integral over P_3 gives the equivalent spectral weight

$$\rho(P_\perp,E)=\frac{1}{(2\pi)^2}R(E^2-P_\perp{}^2) . \tag{4.8}$$

For $\alpha=1$, the Fourier transform has the form of (4.7) with

$$R(\sigma^2)=4\pi\delta(\sigma^2-M^2) . \tag{4.9}$$

It gives the familiar result

$$\rho=\frac{1}{\pi}\delta(E^2-P_\perp{}^2-M^2) . \tag{4.10}$$

To handle $\alpha>1$, use the representation

$$\frac{e^{-Mx}}{x^{1+\epsilon}}=\frac{1}{\Gamma(\epsilon)}\int_M^\infty dM'(M'-M)^{\epsilon-1}\frac{e^{-M'x}}{x} \tag{4.11}$$

valid for $\epsilon>0$. It follows from (4.6), (4.10), and (4.11) that

$$\rho=\frac{1}{2\pi\Gamma(\epsilon)}\frac{\theta(E^2-P_\perp{}^2-M^2)}{(E^2-P_\perp{}^2)^{1/2}}[(E^2-P_\perp{}^2)^{1/2}-M]^{\epsilon-1} . \tag{4.12}$$

This is a positive measure that is consistent with the spectral representation in (4.3) and (4.4).

Now consider $\alpha=0$. Combining

$$e^{-Mx}=-\frac{\partial}{\partial M}\frac{1}{x}e^{-Mx} \tag{4.13}$$

and (4.10) one finds that ρ is the distribution

$$\rho=-\frac{1}{\pi}\frac{\partial}{\partial M}\delta(E^2-P_\perp{}^2-M^2) . \tag{4.14}$$

Since (4.14) is not a positive measure in E at fixed $P_\perp$, it is inconsistent with the spectral representation. Thus the $\alpha=0$, straight-exponential decay is ruled out.

Similarly, the ρ that gives

$$C=\frac{1}{x^\epsilon}e^{-Mx} \text{ with } 0<\epsilon<1 \tag{4.15}$$

is the M derivative of (4.12). As in the $\alpha=0$ case, this ρ is a distribution that is not a positive measure. The conclusion is that values of α less than one are ruled out while the faster decays with $\alpha\geq 1$ are possible.

C. Strong coupling

We have seen that the large-distance behavior of C determined by the spectrum of the transfer matrix near threshold. In this subsection, strong-coupling methods will be applied to the problem.

To precisely identify the property of the spectrum that gives the power correction, substitute (4.10) into (4.3). In performing the E and $P_\perp$ integrations, one finds that the power correction is associated with the $P_\perp{}^2$ term in

$$E(P_\perp)=(P_\perp{}^2+M^2)^{1/2}=M+\frac{1}{2}\frac{P_\perp{}^2}{M}+\cdots . \tag{4.16}$$

Without assuming rotational invariance, we can use strong-coupling methods to show that $E(P_\perp)$ has a $P_\perp{}^2$ term. It is equivalent to the fact that the theory is rough at any finite temperature.[11]

The expansion is in inverse powers of coupling for the plaquettes that have no 3-direction links. It is very similar to Hamiltonian strong coupling[12] except that the coupling and lattice steps for the 3-direction remain finite.

To define the operators, consider the basis

$| \cdots U(l) \cdots \rangle$ in which the group elements on 1-, 2-, and 4-direction links have definite values. The gauge choice has 3-direction links equal to one. Define operators $\hat{U}(l)$, $\hat{R}(l,g_l)$, and $\hat{U}(p)$ by

$$\hat{U}(l) | \cdots U(l) \cdots \rangle = U(l) | \cdots U(l) \cdots \rangle ,$$
$$\hat{U}(p) = \prod_{l \in \partial p} \hat{U}(l) , \tag{4.17}$$
$$\hat{R}(l,g_l) | \cdots U(l) \cdots \rangle = | \cdots g_l^{-1} U(l) \cdots \rangle .$$

The transfer matrix for the Wilson action is then

$$T = BAB \tag{4.18}$$

with

$$A = \prod_{\text{links}} \int dg_l \exp \left[\frac{4}{g^2} \frac{1}{2} \text{Tr}(g_l) \right] \hat{R}(l,g_l) \tag{4.19}$$

and

$$B = \exp \left[\frac{1}{2} \frac{4}{g_s^2} \sum_{\text{plaq}} \tfrac{1}{2} \text{Tr}[\hat{U}(p)] \right] . \tag{4.20}$$

The products and sums do not include 3-direction links or the plaquettes that contain them. The line operator $L(\mathbf{I}_\perp)$ is half the trace of the product of the N_T link operators in the 4-direction at position $\mathbf{I}_\perp$.

The unperturbed basis states for the expansion are the gauge-invariant states that diagonalize A. This is the analogue of the basis used in Hamiltonian strong coupling. The vacuum has all link variables in the trivial representation of the group. Excited states have chains of links in higher representations. The expansion is in powers of $4/g_s^2$.

The operators $L(\mathbf{I}_\perp)$ create states $| \mathbf{I}_\perp \rangle$ that cannot be reconnected to the vacuum in any finite order of the expansion. For sufficiently small $4/g_s^2$, their E values are positive. In lowest order, these states are degenerate and unmixed. However, since N_T is finite, they can be mixed by enough powers of $\ln B$. The states that do not mix have definite momentum

$$| P_\perp \rangle = \sum_{I_\perp} e^{iP_\perp \cdot I_\perp} | \mathbf{I}_\perp \rangle . \tag{4.21}$$

The initial degeneracy of these states is lifted when the expansion connects states with different $\mathbf{I}_\perp$.

If Q is the operator that projects on the space orthogonal to $| P_\perp \rangle$ and D is the perturbation defined by

$$T = BAB = A + D , \tag{4.22}$$

then the series for the shift in the eigenvalues of T is obtained from

$$\Delta\lambda(P_\perp) = \sum_{I_\perp} e^{iP_\perp \cdot I_\perp} \left\langle I_\perp \left| D \frac{1}{1-(QD/\lambda - A)} \right| I_\perp = 0 \right\rangle . \tag{4.23}$$

Since the perturbation can eventually connect the states of different $I_\perp$, the order $P_\perp^2$ term in a momentum expansion of (4.23) does not vanish. It is equivalent to a $P_\perp^2$ term in the expansion

$$E(P_\perp) = M + \tfrac{1}{2} K P_\perp^2 + \cdots . \tag{4.24}$$

This is all we need to know.

These states contribute

$$\delta(E - E(P_\perp)) | \langle P_\perp | L(0) | \text{vac} \rangle |^2 \tag{4.25}$$

to the spectral weight. The corresponding contribution to $C(0,0,I_3)$ is

$$\int d^2 P_\perp e^{-E(P_\perp) I_3} | \langle P_\perp | L(0) | \text{vac} \rangle |^2 . \tag{4.26}$$

For large I_3, the leading term is

$$e^{-MI_3} \int d^2 P_\perp e^{-KP_\perp^2 I_3/2} | \langle P_\perp | L(0) | \text{vac} \rangle |^2$$
$$= \frac{e^{-MI_3}}{I_3} \sqrt{2\pi} \frac{| \langle P_\perp = 0 | L(0) | \text{vac} \rangle |^2}{\sqrt{K}} . \tag{4.27}$$

This result depends upon the $P_\perp^2$ term in E, which is also closely related to roughening.[13] If $E(P_\perp)$ is independent of $P_\perp$, the theory acts one dimensionally and is not rough. With K nonzero, there is roughening and the onset of the restoration of rotational invariance. For any finite N_T, it occurs in a finite order of the expansion. Thus, the finite-temperature theory is always rough. Green has given a detailed discussion.[11]

All of this originates in the ability of the perturbation to move the line from one $\mathbf{I}_\perp$ position to another. These transverse fluctuations are responsible for the power-law correction.

V. CONCLUSIONS

We have discussed the power-law correction to the exponential decay of the line-line correlation function in the finite-temperature, confining phase. Although it does not appear in straightforward series expansions, the correction is quite general and can be derived from the string model, the analogue spin model, or the general properties of the lattice theory at strong coupling. It originates in the transverse fluctuations of the tube of color-electric flux between the quark and antiquark.

ACKNOWLEDGMENTS

It is a pleasure to acknowledge a very helpful conversation with Craig Tracy. We thank Kevin Sparks for a careful reading of the manuscript. This research was supported by the Alfred P. Sloan Foundation and the United States Department of Energy.

[1]Ph. De Forcrand, G. Schierholz, H. Schneider, and M. Teper, Phys. Lett. **160B**, 137 (1985). Other string-model calculations that are very closely related are M. Flensburg and C. Peterson, *ibid.* **153B**, 412 (1985); J. Ambjorn, P. Olesen, and C. Peterson, *ibid.* **142B**, 410 (1984); Nucl. Phys. **B244**, 262 (1984); P. Olesen, *ibid.* **B267**, 539 (1986).

[2]M. Lüscher, K. Symanzik, and P. Weisz, Nucl. Phys. **B173**, 365 (1980); M. Peskin, Phys. Lett. **94B**, 161 (1980); J. Stack

and M. Stone, *ibid.* **100B**, 476 (1981); J. Arvis, *ibid.* **127B**, 106 (1983); O. Alvarez, Phys. Rev. D **24**, 440 (1981); M. Lüscher, Nucl. Phys. **B180** [FS2], 317 (1981).

[3]L. Ornstein and F. Zernike, Z. Phys. **19**, 134 (1918); **27**, 761 (1926); W. Camp and M. Fisher, Phys. Rev. B **6**, 946 (1972).

[4]L. Yaffe and B. Svetitsky, Phys. Rev. D **26**, 963 (1982); B. Svetitsky and L. Yaffe, Nucl. Phys. **B210** [FS6], 423 (1982); F. Green and F. Karsch, *ibid.* **B238**, 297 (1984); M. Ogilvie, Phys. Rev. Lett. **52**, 1369 (1984); A. Gocksch and M. Ogilvie, *ibid.* **54**, 1772 (1985); M. Gross and J. Wheater, Nucl. Phys. **B240** [FS2], 253 (1984); J. Polonyi and K. Szlachanyi, Phys. Lett. **110B**, 395 (1982); L. McLerran and B. Svetitsky, Phys. Rev. D **24**, 450 (1981); Phys. Lett. **98B**, 195 (1981); N. Weiss, Phys. Rev. D **24**, 475 (1981); **25**, 2667 (1982); N. Parga, Phys. Lett. **107B**, 442 (1981); H. Yoneyama, Pisa IFUP Report No. TH83/21, 1983 (unpublished); M. Gross, Phys. Lett. **132B**, 125 (1983); J. Drouffe, J. Jurkiewicz, and A. Krzywicki, Phys. Rev. D **29**, 2982 (1984); E. Gava, R. Jengo, and C. Omero, Nucl. Phys. **B170** [FS1], 445 (1980); L. Susskind, Phys. Rev. D **20**, 2610 (1979); A. Polyakov, Phys. Lett. **72B**, 477 (1978); T. Banks and A. Ukawa, Nucl. Phys. **B225** [FS9], 145 (1983); J. Bartholomew, D. Hochberg, P. Damgaard, and M. Gross, Phys. Lett. **133B**, 218 (1983).

[5]W. Camp, Phys. Rev. B **6**, 960 (1972).

[6]J. Bricmont and J. Fröhlich, Nucl. Phys. **B251** [FS13], 517 (1985); Stack and Stone (Ref. 2).

[7]String-model papers are included in Refs. 1, 2, and 13. In addition, R. Pisarski and O. Alvarez, Phys. Rev. D **26**, 3735 (1982).

[8]Lüscher (Ref. 2).

[9]For example, J. Bjorken and S. Drell, *Relativistic Quantum Fields* (McGraw-Hill, New York, 1965).

[10]J. Kogut, Rev. Mod. Phys. **51**, 659 (1979); E. Seiler, *Gauge Theories as a Problem of Constructive Quantum Field Theory and Statistical Mechanics* (Springer, Berlin, 1982); C. Borgs and E. Seiler, Commun. Math. Phys. **91**, 329 (1983); J. Drouffe and C. Itzykson, Phys. Rep. **38C**, 133 (1978).

[11]F. Green, Nucl. Phys. **B215** [FS7], 83 (1983); **B215** [FS7], 349 (1983).

[12]Kogut (Ref. 10); Drouffe and Itzykson (Ref. 10); J. Kogut, D. Sinclair, and L. Susskind, Nucl. Phys. **B114**, 199 (1976).

[13]J. Kogut, D. Sinclair, R. Pearson, J. Richardson, and J. Shigemitsu, Phys. Rev. D **23**, 2945 (1981); Lüscher (Ref. 2); M. Lüscher, G. Münster, and P. Weisz, Nucl. Phys. **B180** [FS2], 1 (1981); G. Münster and P. Weisz, *ibid.* **B180** [FS2], 13 (1981); **B180** [FS2], 330 (1981); C. Itzykson, M. Peskin, and J. Zuber, Phys. Lett. 95B, 259 (1980); J. Drouffe, J. Luck, and J. Zuber, J. Phys. (Paris) **42**, 653 (1981); A. Hasenfratz, E. Hasenfratz, and P. Hasenfratz, Nucl. Phys. **B180** [FS2], 353 (1981); J. Drouffe and J. Zuber, *ibid.* **B180** [FS2], 253 (1981); **B180** [FS2] 264 (1981).