ERRATUM

Resistance Fluctuations in Thin Bi Wires and Films. D. E. BEUTLER, T. L. MEISENHEIMER, and N. GIORDANO [Phys. Rev. Lett. **58**, 1240 (1987)].

The filled circles in Fig. 2 are data for the wire sample and correspond to the right-hand scale, while the open circles are for the thin-film sample and correspond to the left-hand scale. The terms "open" and "filled" were mistakenly interchanged in the printed version of the caption.

Perturbative Stability of Smooth Strings. ROBERT D. PISARSKI [Phys. Rev. Lett. **58**, 1300 (1987)].

The correct condition for the perturbative stability of smooth strings is that the number of dimensions be positive and less than or equal to thirteen.

Originally, the conformal anomaly for the massless mode was properly treated, but that for the massive mode was overlooked. The inclusion of all terms that arise when a regulator is introduced, then removed by renormalization, adds $(\mathrm{sgn}d)(+1-P/6)$ to the previous result for $\Delta^{-1}(\rho,\rho)$. Equation (7) should read

$$\Delta^{-1}(\rho,\rho)\sim \mathrm{sgn}d\left(\frac{13-d}{3d}P+1-2\frac{\ln P}{P}+\cdots\right), \quad (7')$$

as $P\to\infty$. The characteristic equation for three of the four eigenvalues of Δ^{-1} is then

$$y^3-(cP+1)y^2+\tfrac{1}{2}\ln^2(P)y-\frac{1}{4P}\ln^3(P)=0, \quad (10')$$

up to terms that are negligible at large P. The constant $c=(13-d)/3d$.

When $c\neq 0$, the solutions to Eq. (10′) remain those of Eq. (11), and so the theory is perturbatively stable for $0<d<13$, and unstable for $d>13$.

In thirteen dimensions, c vanishes, and the solutions to Eq. (10′) are

$$y\sim \ln(P)/2P+\cdots,$$
$$y\sim\frac{1}{2}\left[1\pm i\sqrt{2}\ln P\left(1-\frac{1}{4\ln^2 P}\right)\right]+\cdots. \quad (12')$$

Since the real part of each eigenvalue is positive, the theory is perturbatively stable when $d=13$.

At large d, including terms due to the conformal anomaly, about zero momentum the massive mode does not contribute to $\Delta^{-1}(\rho,\rho)$ until $\sim P^2$ (remember P is momentum squared). Hence as $P\to 0$, $\Delta^{-1}(\rho,\rho)\sim -P/6+O(P^2)$. This is precisely the result expected if the infrared limit for $\Delta^{-1}(\rho,\rho)$ is given by the usual Liouville action to $\sim\tilde{\rho}^2$, $\sim(26-d)/6dP\sim -P/6$ at large d. This agrees with Förster, Polyakov, and David, whom I mistakenly contradicted.

A simple picture emerges for the effective action which describes fluctuations in $\tilde{\rho}$,

$$\hat{S}(\rho)=S_{\mathrm{eff}}(\rho,-1\lambda_0\delta^{ab})-S_{\mathrm{eff}}(\rho_0,-i\lambda_0\delta^{ab}).$$

About zero momentum up to terms $\sim P^2$, $\hat{S}(\rho)$ is the standard Liouville action, $\sim(26-d)P$. At large momentum up to corrections ~ 1, $\hat{S}(\rho)$ is again a Liouville-type action, but now it is proportional to $(26-2d)P$. The difference arises because to $\sim P$ about zero momentum, only the massless mode contributes to $\hat{S}(\rho)$, while to leading order at large momentum, the massless and massive modes contribute equally.

The significance of these Liouville actions must be qualified. If mixing between the $\tilde{\rho}$ and $\tilde{\lambda}^{ab}$ fields were ignored, $\hat{S}(\rho)$ would produce logarithmic divergences at both large and short distances, beyond the usual ultraviolet renormalizations of the couplings. Since the full action involves terms $\sim\tilde{\rho}\tilde{\lambda}^{ab}$, Eq. (5), this mixing cannot be ignored. For example, consider the two-point function of $\tilde{\rho}$, $\Delta(\rho,\rho)$. If only $\hat{S}(\rho)$ mattered, $\Delta(\rho,\rho)\sim 1/P$ at both large and small P. Because of the mixing, though, at large P and for any d, $\Delta(\rho,\rho)\sim 1/P\ln^2 P$; e.g., this renders $\langle\tilde{\rho}^2\rangle$ ultraviolet finite. About zero momentum, $\Delta(\rho,\rho)$ can be computed (with apologies) at large d; as asserted before, usually $\Delta(\rho,\rho)$ is finite at $P=0$, with the $\tilde{\rho}$ and $\tilde{\lambda}^{ab}$ fields acquiring a common mass gap.

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