correction for the thermal expansion and neglect the tetragonal distortion of the Rh_2FeSn lattice. [This distortion can influence the values of the coefficients of formula (1).] It is nevertheless clear that the model considered above accounts at least qualitatively for the character of the anomaly of $h(T)$ for Sn at $T/T_c<0.7$.

It is easy to see why formula (3) does not agree with the experimental data at $T/T_c>0.7$. In the critical region of temperatures, the molecular-field model is certainly incorrect: in real ferromagnets, the behavior of the magnetization near T_c does not follow the Brillouin functions. The critical exponent β for the magnetization as given by the molecular-field model is 0.5, whereas the contemporary theoretical and experimental values of β for three-dimensional ferromagnets are close to 0.35–0.38. According to the universality principle, in the critical region the mean values of the local magnetization for the Fe and Rh atoms should have one and the same functional dependence on the temperature (regardless of their behavior at temperatures far from T_c). According to (2), such a temperature dependence should correspond also to the hyperfine field for the Sn atoms. This agrees well with our experimental data: at $T/T_c>0.8$ the $h(T)$ relations for Fe and Sn are the same but do not follow a Brillouin function.

Thus, the characteristic anomalous form of $h(T)$ for Sn in Rh_2FeSn can be fully explained by recognizing that the moments of Fe and Rh have identical temperature dependences near T_c, but substantially different ones in the range $0.4<T/T_c<0.8$. It appears that the temperature anomaly we obtained earlier[14] for the hyperfine field of impurity Sn atoms in the ferromagnetic matrix FeRh is explained in exactly the same way.

It is obvious that similar temperature anomalies of $H(T)$ should be observed also in other systems containing magnetic atoms with substantially different temperature dependences of the local magnetization. A well known example of systems in which the temperature dependences of the magnetization in nonequivalent lattice sites can be substantially different are certain ferrimagnetic compounds. Attention should be called in this connection to the result of Lyubutin *et al.*[15] who obtained a large $H(T)$ anomaly for impurity Sn atoms in the ferrite $MnFe_2O_4$. We note that when the magnetic hyperfine interaction in metallic and nonmetallic systems is considered it is customary to use entirely different concepts and approximations. In particular, no method has been found for representing the hyperfine fields in nonmetals with the aid of universal expressions similar to formula (1), so that the model considered above can not be used directly for an interpretation of the data obtained for nonmetallic ferrimagnets. Nonetheless, it is easy to note that there is an undisputed formal analogy between the model considered by us in the present article and the data interpretation contained in the paper of Lyubutin *et al.*[15]

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Kramers-Wannier transformation for systems with $Z(n)$ symmetry

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Duality relations analogous to the Kramers-Wannier symmetry of the two-dimensional Ising model are established for spin and gauge systems with internal symmetry $Z(n)$.

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Kramers and Wannier[1] drew attention to the fact that the two-dimensional Ising model possesses an exact symmetry that links the high- and low-temperature phases of this model. It is found that the model admits equivalent descriptions, in terms of a spin variable σ (the order parameter) defined on the lattice, and in

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terms of a dual variable μ (the disorder parameter[12]), which is defined on the dual lattice and is also a spin variable, $\mu = \pm 1$. The description of the low temperature phase of the model in terms of σ coincides with the description of the high-temperature phase in terms of μ, and vice versa.

Generalization of the Kramers-Wannier (KW) symmetry to other models is possible in two directions. The first is the generalization to the entire class of models with the same internal symmetry (the symmetry $Z(2)$[1)]) as the Ising model; these can be called generalized gauge Ising models (see Ref. 3).

We shall consider a simple cubic lattice in d-dimensional space. Its dual lattice is also simple cubic, with the sites situated at the centers of the cells of the original lattice. The lattice contains elements of different dimensionality $0 \leq q \leq d$ (for $q = 0$ these are the sites of the lattice, for $q = 1$ its bonds, for $q = 2$ its faces, etc.). The elements of the original lattice that have dimensionality q can be brought into dual correspondence with the $(d-q)$-dimensional elements of the dual lattice. For example, if $d = 3$, the sites of the original lattice are the centers of the cells of the dual lattice, and the bonds pass through the face centers of this lattice.

The generalized gauge Ising models are formulated in the following way: a "generalized gauge field"[2)] A (which is a spin variable $A = \pm 1$) is defined on the lattice elements of dimensionality $q - 1$, and "generalized intensities" F are specified on the q-dimensional elements as the product of the quantities A corresponding to all the $(q-1)$-dimensional elements bounding the given q-dimensional element. A "generalized gauge transformation" is defined by specifying a field σ on the $(q-2)$-dimensional elements of the lattice and multiplying the field A on a given $(q-1)$-dimensional element by the quantities σ specified on all the $(q-2)$-dimensional elements surrounding it. After the gauge transformation, each σ appears twice (in the absence of an overall boundary) in the product defining the intensity F; therefore, the field F is invariant under generalized gauge transformations.

The KW symmetry is formulated for these models in the following way: a model with fields F on the q-dimensional elements of the lattice is dual to a model with fields $\tilde{F}$ defined on the corresponding $(d-q)$-dimensional elements of the dual lattice. For example, the three-dimensional Ising model is dual to its gauge variant, and on a four-dimensional lattice the gauge Ising model is self-dual. Relationships of this type can be established for various mixed models: e.g., the model with a spin field ($q = 1$) interacting with the corresponding gauge field ($q = 2$) is self-dual on a three-dimensional lattice, since the spin field is dual to the gauge field, and vice versa. A review of the topics mentioned here can be found, e.g., in Ref. 4.

The second direction in which the KW symmetry can be generalized is the analysis of systems with other commutative symmetry groups. It is found that the dual systems do not always have the same symmetry group as the original systems. For example, models with internal symmetry $U(1)$ are dual to models with symmetry Z (Z is the group of integers under addition).[5-7]

In the present paper it will be demonstrated that models with internal symmetry $Z(n)$ are dual in the sense mentioned above to models with the same symmetry $Z(n)$, and, for a special choice of the form of the interaction between the spins, the duality is found to be exact, i.e., it can be expressed in the form of a transformation of the temperature, as is the case in the Ising models.

For simplicity, first we shall consider only a spin system on a two-dimensional square lattice, with the global symmetry $Z(n)$. We shall label the lattice sites by the variable $x = \{x_1, x_2\}$, where x_1 and x_2 are integers. The dual lattice is obtained by shifting the origin to $\{\frac{1}{2}, \frac{1}{2}\}$, and its sites are labeled by the variable $\tilde{x} = \{\tilde{x}_1, \tilde{x}_2\}$, where x_1 and x_2 are half-integers. In addition, we introduce two "unit vectors" $\Delta x\mu$, $\mu = 1, 2$, such that $\Delta x_1 = \{1, 0\}$ and $\Delta x_2 = \{0, 1\}$, and their dual pair of "vectors"

$$\widetilde{\Delta x}_\nu = \sum_{\mu=1,2} \varepsilon_{\mu\nu} \Delta x_\mu,$$

where $\varepsilon_{\mu\nu}$ is the antisymmetrizer $\varepsilon_{\mu\nu} = -\varepsilon_{\mu\nu}$, $\varepsilon_{12} = 1$. A bond terminating at the sites x and $x + \Delta x_\mu$ in the original lattice is labeled by the pair x, μ, and a bond of the dual lattice with its ends at $\tilde{x}$ and $\tilde{x} + \widetilde{\Delta x}_\mu$ by the pair $\tilde{x}, \mu$. We note that the bond x, μ is dual to the bond $\tilde{x}, \mu$ if $\tilde{x} = x + \{\frac{1}{2}, \frac{1}{2}\}$.

The elements of the group $Z(n)$, specified at the lattice sites, will be represented by the numbers $\exp(i\Phi_x)$, where $\Phi_x = 2\pi k_x / n$, $k = 0, 1, \ldots, n-1$; here the group multiplication coincides with numerical multiplication. We shall denote the configuration of the system by $\{\Phi_x\}$, and we must not, of course, distinguish between configurations that differ by $2\pi m_x$ if the m_x are integers. The partition function of the model has the form

$$Z(T) = \sum_{\{\Phi_x\}} \exp[-U(\{\Phi_x\}, T)], \tag{1}$$

where T is the temperature parameter. As the functional $U(\{\Phi_x\}, T)$ we take

$$U(\{\Phi_x\}, T) = -\sum_{x,\mu} \ln \sum_{m_{x,\mu}=-\infty}^{\infty} \exp\left[-\frac{1}{2T}(\Phi_x - \Phi_{x+\Delta x_\mu} - 2\pi m_{x,\mu})^2\right]. \tag{2}$$

This complicated and special form of U has been chosen so that the KW transformation does not change the functional form of the interaction energy and reduces to a transformation of the parameter T.[3)] In view of the fact that, for $T \to 0$ and small $\Phi_x - \Phi_{x+\Delta x\mu}$,

$$U(\{\Phi_x\}, T) \to \frac{1}{2T} \sum_{x,\mu} (\Phi_x - \Phi_{x+\Delta x_\mu})^2, \tag{3}$$

we shall call T simply the temperature.

Using (2), we obtain

$$Z(T) = \sum_{\{\Phi_x\}} \sum_{\{m_{x,\mu}\}} \exp\left[-\frac{1}{2T} \sum_{x,\mu} (\Phi_x - \Phi_{x+\Delta x_\mu} - 2\pi m_{x,\mu})^2\right], \tag{4}$$

where $\{m_{x,\mu}\}$ denotes a configuration of integers defined on the bonds of the lattice. If on the bonds we define the increments $\theta_{x,\mu}=\Phi_{x+\Delta x_\mu}-\Phi_x$, the summation over Φ_x can be replaced by a summation over Φ_{x_0} (where x_0 is a certain reference lattice site) and over the configurations $\{\theta_{x,\mu}\}$ of increments. Here, of course, the following condition must be fulfilled:

$$(\mathrm{Rot}\,\theta)_{\tilde{x}}=\theta_{x,1}+\theta_{x+\Delta x_1,2}-\theta_{x+\Delta x_2,1}-\theta_{x,2}=2\pi l_{\tilde{x}}, \tag{5}$$

which takes into account that the circulation of the "vector" $\theta_{x,\mu}$ around a cell of the lattice leads to the starting group element. In formula (5) the lattice "curl" $(\mathrm{Rot}\,\theta)_{\tilde{x}}$ and the intergers $l_{\tilde{x}}$ are defined on the sites of the dual lattice.

It is easy to verify that [4)]

$$\sum_{p=0}^{n-1} e^{ip\Phi}=\begin{cases} n, & \text{if } e^{i\Phi}=1 \\ 0, & \text{if } e^{i\Phi}\in Z(n),\quad e^{i\Phi}\neq 1 \end{cases}. \tag{6}$$

Therefore, the summation in (4) can be taken over all $\{\theta_{x,\mu}\}$ if the factor

$$\prod_{\tilde{x}}\frac{1}{n}\sum_{p_{\tilde{x}}=0}^{n-1}\exp\{ip_{\tilde{x}}(\mathrm{Rot}\,\theta)_{\tilde{x}}\} \tag{7}$$

is introduced into the summation. After this, (4) can be rewritten in the form

$$Z(T)=\sum_{\{\theta_{x,\mu}\}}\sum_{\{m_{x,\mu}\}}\sum_{\{p_{\tilde{x}}\}}\exp\Big\{-\frac{1}{2T}\sum_{x,\mu}(\theta_{x,\mu}+2\pi m_{x,\mu})^2 + i\sum_{x,\mu}\theta_{x,\mu}(p_{\tilde{x}}-p_{\tilde{x}+\tilde{\Delta x}_\mu})\Big\}. \tag{8}$$

Here $\{p_{\tilde{x}}\}$ denotes a configuration of the numbers $p_{\tilde{x}}$. Now the summation over $m_{x,\mu}$ can be performed explicitly and we obtain (we have introduced the notation $\theta_{x,\mu}+2\pi m_{x,\mu}=(2\pi/n)k_{x,\mu}$)

$$Z(T)=\sum_{\{p_{\tilde{x}}\}}\sum_{\{k_{x,\mu}\}}\exp\Big\{-\frac{2\pi^2}{n^2T}\sum_{x,\mu}k_{x,\mu}^2+\frac{2\pi i}{n}\sum_{x,\mu}k_{x,\mu}(p_{\tilde{x}}-p_{\tilde{x}+\tilde{\Delta x}_\mu})\Big\}. \tag{9}$$

Using the well known equality

$$\sqrt{\frac{\alpha}{\pi}}\sum_{m=-\infty}^{\infty}\exp(-\alpha m^2+im\Phi)=\sum_{n=-\infty}^{\infty}\exp\Big[-\frac{1}{4\alpha}(\Phi-2\pi n)^2\Big], \tag{10}$$

we obtain

$$Z(T)\sim\sum_{\{\chi_{\tilde{x}}\}}\sum_{\{l_{\tilde{x},\mu}\}}\exp\Big\{-\frac{1}{2T^*}\sum_{\tilde{x},\mu}(\chi_{\tilde{x}}-\chi_{\tilde{x}+\tilde{\Delta x}_\mu}-2\pi l_{\tilde{x},\mu})^2\Big\}, \tag{11}$$

where $\{l_{\tilde{x},\mu}\}$ denotes a configuration of integers on the bonds of the dual lattice, the notation $\chi_{\tilde{x}}=2\pi p_{\tilde{x}}/n$ has been introduced, and

$$TT^*=4\pi^2/n^2. \tag{12}$$

Finally, it is pertinent to make the following remark. In the above computations we omitted the summation over the variable Φ_{x_0} (which reduces to multiplication by n) and implicitly added an "extra" summation over χ at one of the sites of the dual lattice. To understand the latter, we imagine that the lattice is closed on a torus. If we now pick out some lattice site, the condition curl $\theta=0$ at all the other sites will ensure that the "curl" at this site is equal to zero, so that one summation over p in (7) is superfluous. Because of the $Z(n)$ symmetry of the dual system, this summation leads only to multiplication by n.

The duality relations for "generalized gauged systems" with $Z(n)$ symmetry are derived analogously. We note first that the generalized gauge invariance implies that configurations differing by a gauge transformation must be regarded as identical. Therefore, for systems with generalized gauge symmetry the summation over the "potentials" A can be replaced by a summation over the "intensities" F with the constraint that the lattice "curl" of the field F be equal to zero. This curl is defined on the $(q+1)$-dimensional elements of the original lattice, or, which is equivalent, on the $(d-q-1)$-dimensional elements of the dual lattice that are dual to them. The constraint can be fixed by introducing expressions of the form (7) into the summation, where the numbers p are defined on the $(d-q-1)$-dimensional elements of the dual lattice. This substitution quickly leads to the required duality relations.

We note that for $n=2$ the model (1) conincides with the usual Ising model. Indeed, it is easy to verify that when $n=2$ and when the relation between K and T (K is the inverse temperature of the Ising model) is implicitly determined by the formula

$$e^{-2K}=\sum_{m=-\infty}^{\infty}\exp\Big[-\frac{\pi^2}{2T}(2m+1)^2\Big]\Big/\sum_{m=-\infty}^{\infty}\exp\Big[-\frac{2\pi^2}{T}m^2\Big], \tag{13}$$

the model (1) is transformed into the Ising model. In this case the replacement $T\rightarrow T^*$ corresponds to the replacement $K\rightarrow K^*$, where $\sinh 2K\sinh 2K^*=1$.

The KW symmetry of the two-dimensional Ising model enables us to establish the position of the phase-transition point (from the condition $K=K^*$). For $n\geq 3$, evidently, the situation is different: there are two phase-transition points $T_c^{(1)}$ and $T_c^{(2)}$, with $T_c^{(1)}T_c^{(2)}=4\pi^2/n^2$. The postulated phase diagram in the (T,n)-plane is depicted in the Figure. There is an upper and a lower phase, which transform into each other the KW transformation (the shaded regions in the diagram). In the lower phase the symmetry is broken and the system fluctuates about one of the elements of the group $Z(n)$, while in the upper phase the symmetry of the dual system is broken. Between them is an intermediate region $T_c^{(1)}<T<T_c^{(2)}$. It is possible that in this region the system has the properties of the Berezinskiĭ phase[5,6] of the XY model, i.e., the symmetry is not broken (there is no mean field) but there is transverse stiffness and the correlation functions fall off at large distances in a power-law manner with exponents that

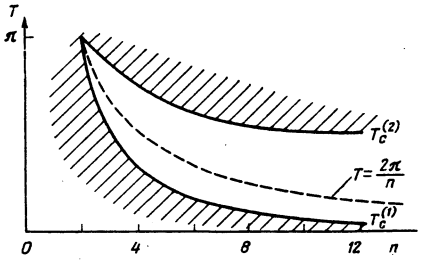

FIG. 1.

depend continuously on the temperature. Then the phase transition at $T=T_c^{(2)}$ corresponds to the vanishing of the transverse stiffness and is analogous to the phase transition in the XY model.

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[1)]The group $Z(n)$ can be defined, e.g., as the set of integers 0, 1, . . . , $n-1$, the group multiplication coinciding with addition modulo n.

[2)]The fields under consideration are a certain formal generalization of a gauge field. A geometrical interpretation of these fields is, however, not known. For $q=2$ the generalized gauge field coincides with the usual one with gauge symmetry $Z(2)$. We note that gauge models with any commutative symmetry can be generalized in a similar formal manner.

[3)]We note that, for a different choice of the interaction energy, all the results cited below remain valid but the KW transformation leads to a change of the functional form of $U(\{\Phi_x\})$.

[4)]The formula (6) is a particular case of the general relation.

$$\sum_{\nu}\chi_{\nu}(G)=\begin{cases} n, & G=I \\ 0, & G\neq I \end{cases}.$$

Here ν labels the irreducible representations of the point group G, n is its order, and $\chi_\nu(G)$ are the characters of these representations. This relation and the analogous formula for continuous groups makes it possible to use the method presented to obtain KW relations for systems with any commutative group.

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Hyperfine magnetic fields at the nuclei ^{57}Fe and ^{119}Sn in the alloy $Fe_{48}Rh_{52}$ under pressure

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The Mössbauer effect is use to measure the pressure dependences of the hyperfine magnetic fields H and of the isomeric shifts ε at the nuclei ^{57}Fe and ^{119}Sn in the alloy $Fe_{48}Rh_{52}$ with ~ 1 at.% Sn as an impurity. Pressure causes ε to decrease, and this corresponds to an increase (for ^{57}Fe) or a decrease (for ^{119}Sn) of the density of the s electrons at the nuclei. In the ferromagnetic (FM) state of the alloy, at 398 K, $\Delta H/H\Delta p = (-2.8\pm0.2)\times10^{-3}$ kbar^{-1} for ^{57}Fe and $(-4.8\pm0.8)\times10^{-3}$ kbar^{-1} for ^{119}Sn; in the antiferromagnetic state (AMF) at 78 K, $\Delta H/H\Delta p\approx 0$ for ^{57}Fe and $\Delta H/H\Delta p = (-6.2\pm1.0)\times10^{-3}$ kbar^{-1} for ^{119}Sn. The results are attributed to the strong dependence of the magnetization of the alloy matrix on the pressure for ^{57}Fe in the FM state and to the absence of "local" polarization of the s-like collectivized electrons and to the pressure dependence of the magnetic moments of the Fe ions in the AFM state. The causes of the different effects of pressure on the magnetic moments of the Fe ions in the FM and AFM states are discussed. The results for ^{119}Sn in the FM and AFM states agree with a previously proposed model [A. E. Balabanov, N. N. Delyagin, *et al.*, Sov. Phys. JETP **27**, 752 (1968) and elsewhere; I. N. Nikolaev and V. P. Potapov, *ibid.* **45**, 840 (1977)] of the hyperfine fields at the Sn impurity atoms in magnetic matrices. An estimate is obtained of the radial dependence of the hyperfine field at the ^{119}Sn nuclei for the AFM state, namely, $H(r)$ varies more strongly than r^{-9}.

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This is a logical continuation of a number of preceding studies[1-6] of the influence of pressure on hyperfine interactions in magnets. The purpose of these studies was to attempt to explain the mechanism whereby hyperfine magnetic fields are produced at nuclei of atoms in magnetic matrices and, in particular, determine the role of spin polarization of collectivized electrons in the onset of magnetic order. In alloys of the Fe_xRh_{1-x} system, when the composition or the temperature is changed, a transformation from the ferromagnetic (FM) into the antiferromagnetic (AFM) state is observed, the parameters of the crystal lattice change jumpwise by 0.3%, and the structure remains cubic. In this case there is a rare opportunity of tracing, in samples having the same composition, the influence of pressure (interatomic distance) on the hyperfine magnetic fields at the nuclei of the matrix (^{57}Fe) and impurity (^{119}Sn) atoms in two magnetic states, FM and AFM, which are produced by varying the temperature an which differ in the orientation of the magnetic moments of the iron atoms, and hence in the polarization of the conduction electrons. For the Fe_xRh_{1-x} alloys, detailed studies were made of the magnetic fields at the nuclei ^{57}Fe (Ref. 7) and ^{119}Sn (Ref. 8), of the distributions of the magnetic moments and of the spin density,[9] and also of the influence of pressure on the temperatures of the

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